Two-body and Three-body Decays of Charginos in One-loop Order in MSSM

FUJIMOTO Junpei a, ISHIKAWA Tadashi a, JIMBO Masato b, KON Tadashi c, KURIHARA Yoshimasa a, KURODA Masaaki d

\( ^{a} \) KEK, Oho, Tsukuba, Ibaraki 305-0801 Japan
\( ^{b} \) Tokyo Management College, Ichikawa, Chiba 272-0001, Japan
\( ^{c} \) Seikei University, Musashino, Tokyo 180-8633, Japan
\( ^{d} \) Meiji Gakuin University, Totsuka, Yokohama 244-8539, Japan

Abstract
We present the two-body and three-body decay widths of charginos in one-loop order in the MSSM.

1 Introduction
The concept of supersymmetry (SUSY) \(^{1}\) \(^{2}\) is considered the most promising extension of the Standard Model (SM) \(^{3}\) of particle physics. Among the supersymmetric theories, the minimal supersymmetric standard model (MSSM) \(^{2}\) is the most elaborated and well studied framework of SUSY. The existence of many new supersymmetric particles in any SUSY model makes it very complicated and tedious to compute even a simple two-body decay width exactly. To overcome this problem, the Minami-tateya group of KEK has constructed a computational system \textit{GRACE/SUSY} \(^{4} \) \(^{5}\) which automatically creates, for a given process, all the Feynman diagrams and compute the Feynman amplitudes and subsequently the cross section or the decay width itself at tree level in the MSSM.

With the increase of more precise experimental data, in particular, in future colliders, it becomes apparent that the inclusion of at least one-loop

\(^{1}\text{Dedicated to Dr. J. Kodaira, our friend and colleague, who passed away in Sept. 2006}\)
radiative corrections is necessary for calculations of cross sections of SUSY processes. We have, therefore, extended GRACE/SUSY in order to incorporate radiative corrections in one-loop order in the system. The new system, called GRACE/SUSY-loop, is constructed based on the same philosophy as for GRACE-loop [6] developed by the Minami-tateya group, which is the automatic computation system for the SM processes including one-loop corrections. Some of the results computed by this system have already been presented in several workshops and symposiums [7] [8] [9].

In this paper we present the result of the thorough investigation of the two-body and three-body decays of charginos in one-loop order in the MSSM computed by the automatic computing system GRACE/SUSY-loop. We also show the cross section of the chargino pair production in $e^+e^−$ annihilation and the subsequent decays by combining the production cross section and the decay rates.

The paper is organized as follows. In section two, the renormalization scheme used in GRACE/SUSY-loop is explained. The features of the GRACE/SUSY-loop system are briefly given in section three. The numerical results are presented in section four, and several comments on the numerical results are given in section five.

2 Renormalization scheme

In this section we explain briefly the renormalization scheme adopted in GRACE/SUSY-loop. Our approach is a straightforward extension of the on-shell renormalization in the SM used in GRACE [6]. The essential part of the scheme has been given in [10].

The Lagrangian of the MSSM has been given in [11] [12] [13]. In terms of superfields, it is given as (see [11] for detail)

\[
\mathcal{L} = \int d^2\theta \frac{1}{4}[2Tr(W^2) + WW + 2Tr(W_sW_s)] + h.c.
+ \int d^2\theta d^2\Phi^\dagger \exp[2(g_s^α V^α + g_s \frac{Y_q}{2} V)]\Phi^\dagger
+ \int d^2\theta d^2\Phi_{\ell} \exp(g'Y_{\ell} V)\Phi_{\ell}
+ \int d^2\theta d^2\Phi_{q} \exp[2(g_s^α V^α + g_s \frac{Y_q}{2} V + g_s \lambda^α \frac{V^α}{2})]\Phi_{q}
+ \int d^2\theta d^2\Phi_{u} \exp(g'Y_{u} V - g_s \lambda^α \frac{V^α}{2})\Phi_{u}
\]
+ \int d^2 \vartheta \bar{\Theta} d \Phi_H \exp(g' Y_d V - g_s \lambda^\alpha V^\alpha) \Phi_d
+ \int d^2 \vartheta \bar{\Theta} d \Phi_H \exp[2(g T^a V^a + g' Y_d H_1 V)] \Phi_H
+ \int d^2 \vartheta \bar{\Theta} d \Phi_H \exp[2(g T^a V^a + g' Y_d H_2 V)] \Phi_H
+ \sqrt{2} m_e \int d^2 \vartheta \Phi_H \Phi_e + h.c.
- \sqrt{2} m_u \int d^2 \vartheta \Phi_H \Phi_u + h.c.
+ \sqrt{2} m_d \int d^2 \vartheta \Phi_H \Phi_d + h.c.
- \mu \int d^2 \vartheta \Phi_H \Phi_H + h.c.
+ \mathcal{L}_{soft}
+ \mathcal{L}_{gf} + \mathcal{L}_{ghost},

(2.1)

where \( W, W, W \) are superfield strengths corresponding to the \( SU(2)_L, U(1) \) and \( SU(3)_c \) gauge-superfields, \( V, V, V \), respectively. The component fields belonging to each superfield are given as follows

\[
\Phi_H = \begin{pmatrix} H_1^0 & \tilde{H}_1^- \\ \tilde{H}_1^+ & \tilde{H}_1^0 \end{pmatrix},
\Phi_H = \begin{pmatrix} H_2^0 & \tilde{H}_2^- \\ \tilde{H}_2^+ & \tilde{H}_2^0 \end{pmatrix},
\Phi_\ell = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix},
\Phi_e = \begin{pmatrix} e_R \\ \tilde{e}_R \end{pmatrix},
\Phi_q = \begin{pmatrix} u_L \\ d_L \end{pmatrix},
\Phi_u = \begin{pmatrix} u_R \\ \tilde{u}_R \end{pmatrix},
\Phi_d = \begin{pmatrix} d_R \\ \tilde{d}_R \end{pmatrix},
\]

(2.2)

The \( SU(2)_L \) doublet gauge-bosons and gauginos are denoted by \( \tilde{W}_\mu, \tilde{W}_\mu \), while the singlet gauge-boson and gaugino are denoted by \( B_\mu \) and \( \lambda \). The Higgs
and higgsino doublets are denoted by $H_i$ and $\tilde{H}_i$ with $i = 1, 2$, respectively. The soft SUSY breaking terms are expressed as

$$L_{soft} = -\frac{1}{2} M_1 \lambda - \frac{1}{2} M_2 \lambda^2 \lambda - \frac{1}{2} M_3 \tilde{g}^2 \tilde{g}^2 + h.c.$$  

$$-\tilde{m}_1^2 H_1^2 - \tilde{m}_2^2 H_2^2 - (\tilde{m}_{12}^2 H_1 H_2 + h.c.) - \sum \tilde{m}_{f_m}^2 \tilde{f}_m^* \tilde{f}_m$$  

\begin{equation}
-\sqrt{2} m_u A_u H_2 A_L (u_L) + \sqrt{2} m_d A_d H_1 A_L (d_L) + h.c.
\end{equation}

\begin{equation}
\sqrt{2} m_e A_e H_1 A_L (e_L) + h.c. (2.3)
\end{equation}

where the sum in the sfermion mass terms runs for $\tilde{f}_m = \tilde{f}_L$ and $\tilde{f}_R$. The sign convention of our $A_f$ is opposite to the convention used by others.

Sfermion mass eigenstates are denoted by $\tilde{f}_i$ with $i = 1, 2$ which are the mixture of the left-handed ($\tilde{f}_L$) and the right-handed sfermions ($\tilde{f}_R$). The sfermion mass matrix is diagonalized as,

$$\begin{bmatrix}
\cos \theta_f & \sin \theta_f \\
-\sin \theta_f & \cos \theta_f
\end{bmatrix}
\begin{bmatrix}
m_{f_L}^2 \\
m_{f_R}^2
\end{bmatrix}
= \begin{bmatrix}
m_{f_{L1}}^2 & 0 \\
0 & m_{f_{L2}}^2
\end{bmatrix},$$

\begin{equation}
(2.4)
\end{equation}

where

$$m_{f_L}^2 = \tilde{m}_{f_L}^2 + m_{f}^2 + M_2^2 \cos 2\beta (T_{3f} - Q_f s_{W}^2),$$

$$m_{f_R}^2 = \tilde{m}_{f_R}^2 + m_{f}^2 + M_2^2 \cos 2\beta Q_f s_{W}^2,$$

\begin{equation}
m_{f_{LR}}^2 = \begin{cases} 
-m_u (\mu \cot \beta + A_u), & f = u \\
-m_d (\mu \tan \beta + A_d), & f = d, e
\end{cases}
\end{equation}

Their masses and the mixing angles satisfy the relations originating from the SU(2)$_L$ conditions on their left-handed soft SUSY-breaking mass terms, $\tilde{m}_{u_L}^2 = \tilde{m}_{d_L}^2$, $\tilde{m}_{e_L}^2 = \tilde{m}_{\nu_e}^2$ etc. For the third generation of sfermions, for example,

$$\cos^2 \theta_t m_{t_1}^2 + \sin^2 \theta_t m_{t_2}^2 - m_t^2 = \cos^2 \theta_b m_{b_1}^2 + \sin^2 \theta_b m_{b_2}^2 - m_b^2 + M_W^2 \cos 2\beta,$$

$$m_{\nu_\tau}^2 = \cos^2 \theta_{\tau} m_{\tau_1}^2 + \sin^2 \theta_{\tau} m_{\tau_2}^2 - m_{\tau}^2 + M_W^2 \cos 2\beta.$$  

\begin{equation}
(2.6)
\end{equation}
The renormalization constants are introduced as follows.

[Standard Model sector]

\[ \bar{W}_{\mu 0} = Z_{W}^{1/2} \bar{W}_{\mu}, \]
\[ B_{\mu 0} = Z_{B}^{1/2} B_{\mu}, \]
\[ g_{\mu 0} = Z_{g}^{1/2} g_{\mu}, \]

[Neutral gauge bosons]
\[ g_{0} = Z_{g}^{-3/2} g, \]
\[ g'_{0} = Z_{g}'^{-3/2} g', \]
\[ g_{s0} = Z_{g_{s}}^{-3/2} g_{s}, \]

[Neutral gauge couplings]
\[ g_{0} = Z_{g}^{-3/2} g, \]
\[ g'_{0} = Z_{g}'^{-3/2} g', \]
\[ g_{s0} = Z_{g_{s}}^{-3/2} g_{s}, \]

[Fermions]
\[ \Psi_{fL 0} = Z_{f}^{L} \frac{1}{2} \Psi_{fL}, \quad f = u, d, \ldots, \nu, e, \ldots, \]
\[ \Psi_{fR 0} = Z_{f}^{R} \frac{1}{2} \Psi_{fR}, \quad f = u, d, \ldots, e, \ldots, \]
\[ m_{f0} = m_{f} + \delta m_{f}, \quad f = u, d, \ldots, e, \ldots. \] (2.7)

[SUSY sector]

| Higgs bosons | $H_{i 0} = Z_{H_{i}}^{1/2} H_{i}$, $i = 1, 2$, |
|-------------|---------------------------------------------|
|             | $v_{i 0} = Z_{H_{i}}^{1/2} (v_{i} - \delta v_{i})$, $i = 1, 2$, |
|             | $m_{i 0} = Z_{H_{i}}^{-1/2} (m_{i}^{2} + \delta m_{i}^{2})$, $i = 1, 2$, |
|             | $(m_{12}^{2})_{0} = Z_{H_{1}}^{-1/2} Z_{H_{2}}^{-1/2} (m_{12}^{2} + \delta m_{12}^{2})$, |
| sfermions   | $\begin{pmatrix} f_{1} \\ f_{2} \end{pmatrix}_{0} = \begin{pmatrix} Z_{f_{1} f_{1}}^{1/2} & Z_{f_{1} f_{2}}^{1/2} \\ Z_{f_{2} f_{1}}^{1/2} & Z_{f_{2} f_{2}}^{1/2} \end{pmatrix} \begin{pmatrix} f_{1} \\ f_{2} \end{pmatrix}$, $f = u, d, \ldots, e, \ldots$, |
|             | $(\tilde{\nu}_{i})_{0} = Z_{\tilde{\nu}_{i}}^{1/2} \tilde{\nu}_{i}$, $i = e, \mu, \tau$, |
|             | $(m_{f_{i}}^{2})_{0} = m_{f_{i}}^{2} + \delta m_{f_{i}}^{2}$, $f = u, d, \ldots, e, \ldots$, $i = 1, 2$, |
|             | $(m_{\tilde{\nu}_{i}}^{2})_{0} = m_{\tilde{\nu}_{i}}^{2} + \delta m_{\tilde{\nu}_{i}}^{2}$, $i = e, \mu, \tau$, |
|             | $(\theta_{f})_{0} = \theta_{f} + \delta \theta_{f}$, $f = u, d, \ldots, e, \ldots$, |
| Inos        | $\bar{\lambda}_{0} = Z_{\bar{\lambda}}^{1/2} \bar{\lambda}$, |
|             | $\lambda_{0} = Z_{\lambda}^{1/2} \lambda$, |
|             | $\tilde{H}_{i 0} = Z_{H_{i}}^{1/2} \tilde{H}_{i}$, $i = 1, 2$, |
\[ g_0^\alpha = Z_\tilde{g}^{1/2}g^\alpha, \]
\[ g_0 = \mu + \delta \mu, \]
\[ M_{10} = M_1 + \delta M_1, \]
\[ M_{20} = M_2 + \delta M_2, \]
\[ M_{30} = M_3 + \delta M_3, \quad (m_{\tilde{g}_0} = m_{\tilde{g}} + \delta m_{\tilde{g}}), \quad (2.8) \]

where in the Higgs boson part,
\[ m_i^2 = \tilde{m}_i^2 + |\mu|^2, \quad i = 1, 2. \quad (2.9) \]

Summing up, we have \((3 + 7N_G)\) wavefunction renormalization constants, \((3 + 3N_G)\) mass counterterms in the non-SUSY sector, and \((7 + 13N_G)\) wavefunction renormalization constants and \((9 + 10N_G)\) mass, vacuum-expectation-value and mixing-angle counterterms in the SUSY sector.

Several comments are in order. In the Higgs and higgsino sectors, the wavefunction renormalization constants are introduced to each unmixed bare doublet state. The mixing angles in the Higgs, chargino and neutralino sectors are defined as the angles which diagonalize the renormalized mass matrices. Therefore, there appear no bare mixing angles or counterterms for the mixing angles of charginos and neutralinos in our scheme. See [14].

In place of \(\delta m_1^2, \delta m_2^2\) and \(\delta m_{12}^2\), we use the mass counterterm of the CP odd Higgs particle, \(\delta M_A^2\), and two counterterms of the tadpole interactions, \(\delta T_1\) and \(\delta T_2\), which are given by the linear combination of \(\delta m_1^2, \delta m_2^2\) and \(\delta m_{12}^2\):

\[ \delta M_A^2 = s_\beta^2 \delta m_1^2 + c_\beta^2 \delta m_2^2 - 2c_\beta s_\beta \delta m_{12}^2 \frac{M_Z^2}{2}(c_\beta^2 - s_\beta^2)^2\delta Z_H_1 + \delta Z_H_2 \frac{2c_\beta}{c_\beta^2 - s_\beta^2} \delta v_1 + \frac{2s_\beta^2}{c_\beta^2 - s_\beta^2} \delta v_2, \quad (2.10) \]

\[ \delta T_1 = v_1 \delta m_1^2 - m_1^2 \delta v_1 + v_2 \delta m_{12}^2 - m_{12}^2 \delta v_2 \]
\[ + \frac{1}{8}(g^2 + \delta g^2)(v_1^2 - v_2^2) \delta v_1 + \frac{1}{8}(g^2 + g'^2)(v_2^2 - 3v_1^2) \delta v_1 + 2v_1 v_2 \delta v_2, \]
\[ + \frac{1}{8}(g^2 + g'^2)(2v_1^3 - v_1 v_2^2) \delta Z_{H_1} - v_1 v_2^2 \delta Z_{H_2}, \quad (2.11) \]

\[ \delta T_2 = (1 \leftrightarrow 2) \quad \text{in} \quad \delta T_1. \quad (2.12) \]

We introduce the gauge-fixing Lagrangian in terms of the renormalized fields as we have done in GRACE-loop for SM [6]. No renormalization constants are introduced for gauge fixing constants and no counterterm La-
The renormalization of the ghost fields is not necessary in one-loop order.

We use the on-shell renormalization scheme. The renormalization conditions employed in GRACE/SUSY-loop are the following set of conditions. Using these conditions, we can express all the renormalization constants and counterterms in terms of the linear combination of the two-point functions evaluated at some specific renormalization points.

**gauge sector**

We use the on-mass-shell condition for $W$, $Z$ and photon. In addition, we require that the residue of the photon propagator at the pole position is one, and that the $A_\mu-Z_\mu$ transition vanishes for the on-shell photon.

**SM fermions**

We use the same renormalization conditions adopted in GRACE-loop. Namely, we require the on-mass-shell condition and the residue condition that the residue of the fermion propagator at the pole is one.

**Higgs sector**

We impose the on-mass-shell and the residue condition for CP odd Higgs, $A^0$, and the decoupling of $Z_\mu$ and $A^0$ on the mass-shell of $A^0$. For the CP even Higgs, we impose the on-mass-shell condition for the heavier Higgs, $H^0$. The above three conditions together with one of the conditions imposed in the gauge sector determine four renormalization constants $\delta H_1, \delta H_2, \delta v_1$ and $\delta v_2$.

Note that we have not adopted the often erroneously used renormalization condition, $\delta v_1 = \delta v_2$, since this condition violates the gauge invariance [15].

**tadpole terms**

Identical to GRACE-loop, we require that the tadpole terms in the renormalized Lagrangian vanish by itself and the tadpole counterterms cancel the one-loop tadpole contributions.

**chargino sector**

We impose the on-mass-shell condition on both $\tilde{\chi}_1^+$ and $\tilde{\chi}_2^+$. In addition, we impose the residue condition on $\tilde{\chi}_1^+$.

**neutralino sector**

We impose the on-mass-shell and the residue condition on the lightest neutralino, $\tilde{\chi}_1^0$.

**sfermion sector**

We impose the on-mass-shell condition and the residue condition on all the
seven sfermions in each generation. In addition, we impose that there is no induced mixing between physical \( \tilde{f}_1 \) and \( \tilde{f}_2 \). The counterterm for the slepton mixing angle is determined by the SU(2)\(_L\) relation between the bare slepton masses. In the squark sector, there are two counterterms \( \delta \theta_u \) and \( \delta \theta_d \) for each generation. We fix \( \delta \theta_u \) \([16]\) by

\[
\delta \theta_u = \frac{1}{2} \left( \Sigma_{\tilde{u}_1} \tilde{u}_2 (m_{\tilde{u}_1}^2) + \Sigma_{\tilde{u}_1} \tilde{u}_2 (m_{\tilde{u}_2}^2) \right) \frac{m_{\tilde{u}_2}^2 - m_{\tilde{u}_1}^2}{m_{\tilde{u}_2}^2 - m_{\tilde{u}_1}^2},
\]

while \( \delta \theta_d \) is fixed by the SU(2)\(_L\) relation between the bare squark masses.

**Charge renormalization**

The charge (electromagnetic coupling constant) is defined, as in the standard model \([6]\) by the Thomson limit of the photon-electron-electron vertex.

**QCD sector**

We impose the on-mass-shell condition and the residue condition for gluons and gluinos. The counterterm \( \delta Z_{g_s} \) is determined by the minimal subtraction with dimensional reduction \((\overline{DR})\) \([6]\).

The explicit expression of the renormalization constants in terms of two-point functions is given in Appendix A.

A couple of comments are worthwhile at this stage. In addition to the wavefunction renormalization constants introduced in \((2.7)\) and \((2.8)\), we need to introduce the ultraviolet finite external wavefunction renormalization constant \( \delta Z^{ext} \) for each particle for which the residue condition is not imposed on its propagator, namely for \( W^\pm, Z^0, H^0, h^0, H^\pm, \tilde{\chi}_0^\pm, \tilde{\chi}_1^0, \tilde{\chi}_2^0, \tilde{\chi}_3^0, \tilde{\chi}_4^0 \). The expression of \( \delta Z^{ext} \) in terms of two-point functions is given in Appendix B.

The counterterm \( \delta A_f \) of the coupling constant of the soft SUSY-breaking Yukawa interaction among Higgs and sfermions is given by

\[
\delta (m_f A_f) = \frac{1}{2} (\delta m_{f_{j2}}^2 - \delta m_{f_{j1}}^2) \sin 2\theta_f + \delta \theta_f (m_{f_{j2}}^2 - m_{f_{j1}}^2) \cos 2\theta_f
\]

\[
- \begin{cases} 
\delta (m_{f_\mu} \cot \beta) & f = u, c, t \\
\delta (m_{f_\mu} \tan \beta) & f = d, s, b, e, \mu, \tau
\end{cases}
\]

since according to \((2.4)\) \( A_f \) is related to the sfermion masses \( m_{f_{j1}}^2 \) and \( m_{f_{j2}}^2 \) as

\[
m_f A_f = \cos \theta_f \sin \theta_f (m_{f_{j2}}^2 - m_{f_{j1}}^2) - \begin{cases} 
m_{f_\mu} \cot \beta & f = u, c, t \\
m_{f_\mu} \tan \beta & f = d, s, b, e, \mu, \tau
\end{cases}.
\]

The system can easily accommodate different renormalization schemes by re-expressing the renormalization constants and the mass counterterms in
terms of different linear combinations of the two-point functions.

A brief comparison of our scheme with other earlier studies is worthwhile to clarify the difference and also the possible scheme dependence in different schemes. There are plenty of papers on the renormalization schemes of the MSSM, and it is beyond our scope to compare all of them. The earlier works on the MSSM renormalization (see for example, [17, 14, 18]) are naturally concerned with radiative corrections in the Higgs sector. Later, the study of the renormalization of the other sector, gaugino and higgsino sector and sfermion sector followed (see for example, [19, 20, 21, 22]). In [17, 18, 14, 21, 23] the on-shell renormalization is used, while the so-called DR is used in [20].

The renormalization scheme we adopt in GRACE/SUSY-loop is close to the scheme given by [14], but there are a couple of differences. In [14], different from our scheme, the gauge-fixing Lagrangian is introduced in terms of bare fields. Therefore, in their scheme, they need extra renormalization conditions to fix the renormalization constants for gauge-parameters. Another difference of our scheme from the others lies in the renormalization condition imposed on $\delta v_i$. The condition used in [18, 14, 21] $\delta v_1/v_1 = \delta v_2/v_2$ is not preferable on the ground of potentially violating the Ward identity. We use the on-mass-shell condition for $H^0$ in place of the condition on the vacuum expectation value, which leads to our expression of $\delta \tan \beta$ which is also different from others.

3 Features of GRACE/SUSY-loop

We present in this section several features of the current version of the system GRACE/SUSY-loop. Some of them have been given in [7].

In order to check and detect possible errors in the system we have used the technique of the non-linear gauge (NLG) [24]. Since physical results are independent on the NLG parameters, they must vanish in the sum of the Feynman amplitudes for physical processes. The test using the NLG parameters provides us with a more powerful tool for the check of the system than the test using the linear gauge parameters, because each NLG parameter is concerned with many kinds of amplitudes which are not in the same gauge-independent sub-set in the linear gauge [9].

In the actual computation of decay widths and production cross sections
in **GRACE/SUSY-loop**, we use the 't Hooft-Feynman gauge with $\xi_V = 1$. For the consistency check of the computation, we use the vanishing of the ultraviolet divergences and the infrared singularities in the sum of the loop and the soft photon/gluon radiation diagrams, as well as the stability in the sum of the soft photon/gluon radiation and the hard photon/gluon radiation diagrams against the change of the photon/gluon energy cut-off.

The input parameters of **GRACE/SUSY-loop** are

$$e, g_s, M_W, M_Z, M_A^0, \tan \beta, \mu, M_1, M_2, M_3, \phantom{1^2}$$
$$m_u, m_d, m_e, \cdots, \phantom{1^2}$$
$$m_{\tilde{u}_1}, m_{\tilde{u}_2}, m_{\tilde{d}_1}, m_{\tilde{d}_2}, m_{\tilde{e}_1}, m_{\tilde{e}_2}, m_{\tilde{\nu}_e}, \cdots, \phantom{1^2}$$
$$\theta_u, \theta_d, \theta_e, \cdots, \phantom{1^2}$$

Using (2.6), we fix the remaining two masses of the sfermions in each generation.

The coupling constants $A_f$ of the soft SUSY-breaking Yukawa interaction among Higgs and sfermions are not our independent input parameters, since they are expressed in terms of sfermion masses and the mixing angles as (2.15). We don’t use the GUT relations $M_1 = \frac{5}{3} \tan^2 \theta_W M_2$, $M_3 = \frac{g_2}{\tan^2 \theta_W} M_2$ at our energy scale.

In our renormalization scheme, the pole mass of $h^0, H^\pm, \tilde{\chi}_0^0, \tilde{\chi}_0^2, \tilde{\chi}_0^3$ and $\tilde{\chi}_0^4$ is different from its Born value.

We have chosen the input values of the parameters in such a way that they reproduce the the SPA1 pole mass values [25] as close as possible. Since the input values of the SPA1 parameters [25] are defined by the $\overline{DR}$ scheme, it is not possible to reproduce exactly the same values as those proposed by SPA1. Note that the lightest neutralino ($\tilde{\chi}_0^1$) is the lightest SUSY-particle (LSP) both in SPA1 and in our parameter choice.

In the numerical calculation we adopt two numerical sets (A) and (B), which are given in Table 1 and 2, respectively. For the fermion masses which are not listed in Tables 1 and 2, we use the standard value. The (one-loop improved) mass of Higgs particles, charginos and neutralinos in the set (A) and (B) is given in Table 3 and 4, respectively.

The set (A) is almost the same as the SPA1a' parameter set, in which not only the heavier chargino $\tilde{\chi}_1^+$ but also the lighter chargino $\tilde{\chi}_1^+$ can decay into various two bodies because the lighter chargino is heavier than some sleptons. For the set (B), on the other hand, the mass of the lighter chargino $m_{\tilde{\chi}_1^+}$ is smaller than the mass of all sfermions $m_{\tilde{f}}$ as well as the sum $m_W + m_{\tilde{\chi}_1^0}$.
| \( \tan \beta \) | \( \mu \) | \( M_1 \) | \( M_2 \) | \( M_3 \) | \( M_{A^0} \) |
|---|---|---|---|---|---|
| 10.00 | 399.31 | 100.12 | 197.52 | 610 | 424.9 |

| \( m_{\tilde{u}_1} \) | \( m_{\tilde{u}_2} \) | \( m_{\tilde{d}_1} \) | \( m_{\tilde{d}_2} \) | \( m_{\tilde{e}_1} \) | \( m_{\tilde{e}_2} \) | \( m_{\tilde{\mu}_1} \) | \( m_{\tilde{\mu}_2} \) | \( m_{\tilde{\nu}_e} \) | \( \cos \theta_u \) | \( \cos \theta_d \) | \( \cos \theta_{\mu} \) |
|---|---|---|---|---|---|---|---|---|---|---|---|---|
| 545.67 | 563.44 | 545.50 | 569.03 | 125.50 | 190.14 | 172.70 | 2.4 \times 10^{-3} | 0.011 | 1.1 \times 10^{-1} |

| \( m_{\tilde{\tau}_1} \) | \( m_{\tilde{\tau}_2} \) | \( m_{\tilde{b}_1} \) | \( m_{\tilde{b}_2} \) | \( m_{\tilde{t}_1} \) | \( m_{\tilde{t}_2} \) | \( m_{\tilde{\nu}_\tau} \) | \( \cos \theta_t \) | \( \cos \theta_b \) | \( \cos \theta_{\tau} \) |
|---|---|---|---|---|---|---|---|---|---|
| 368.53 | 583.79 | 450.12 | 544.38 | 107.71 | 195.08 | 172.69 | 0.063 | 0.018 | 0.023 |

Table 1: The value of the MSSM input parameters for set (A)

| \( h^0 \) | \( H^0 \) | \( A^0 \) | \( H^\pm \) |
|---|---|---|---|
| 107.12 | 425.30 | 424.90 | 432.75 |

| \( \tilde{\chi}_1^+ \) | \( \tilde{\chi}_2^+ \) | \( \tilde{\chi}_1^0 \) | \( \tilde{\chi}_2^0 \) | \( \tilde{\chi}_3^0 \) | \( \tilde{\chi}_4^0 \) |
|---|---|---|---|---|---|
| 184.2 | 421.2 | 97.75 | 184.62 | 398.30 | 413.39 |

Table 3: The pole mass of Higgs, charginos and neutralinos in set (A)

| \( h^0 \) | \( H^0 \) | \( A^0 \) | \( H^\pm \) |
|---|---|---|---|
| 122.50 | 431.40 | 431.0 | 438.73 |

| \( \tilde{\chi}_1^+ \) | \( \tilde{\chi}_2^+ \) | \( \tilde{\chi}_1^0 \) | \( \tilde{\chi}_2^0 \) | \( \tilde{\chi}_3^0 \) | \( \tilde{\chi}_4^0 \) |
|---|---|---|---|---|---|
| 147.08 | 418.8 | 97.61 | 147.4 | 404.0 | 418.8 |

Table 4: The pole mass of Higgs, charginos and neutralinos in set (B)
This means that the lighter chargino $\tilde{\chi}_1^+$ cannot decay into any two bodies and $\tilde{\chi}_1^+$ has only three body decay modes, $f\bar{f}\tilde{\chi}_1^0$.

### 4 Numerical results

The one-loop electroweak corrections on various two-body decay widths of the lighter $\tilde{\chi}_1^+$ and heavier chargino $\tilde{\chi}_2^+$ for the parameter set (A) are shown in Table 5 and 6, respectively, where we do not display the decay modes with small branching fraction $\text{Br}<0.1\%$. While all decays in Table 5 are the electroweak processes, the process $\tilde{\chi}_2^+ \rightarrow \tilde{b}_1$ in Table 6 gets both the electroweak and QCD corrections through the loop contributions and the photon/gluon emissions. We define $\Gamma \equiv \Gamma_0 + \delta\Gamma$, where $\Gamma_0$ and $\delta\Gamma$ are the improved Born decay width and the one-loop correction, respectively. Note that the improved Born decay width is different from the Born width. We obtain $\Gamma_0$ by replacing the tree-level masses by the one-loop renormalized pole-masses presented in Tables 3 and 4 in the tree amplitudes.

| Decay Mode         | $\Gamma_0$ (GeV) | $\Gamma$ (GeV) | $\delta\Gamma/\Gamma_0$ | Br  |
|--------------------|-------------------|----------------|--------------------------|-----|
| $\tilde{\chi}_1^+ \rightarrow \nu_\tau \tilde{\tau}_1^+$ | $3.91 \times 10^{-2}$ | $3.78 \times 10^{-2}$ | $-3.3\%$ | 50.11% |
| $\tilde{\chi}_1^+ \rightarrow \nu_\mu \tilde{\mu}_1^+$ | $1.33 \times 10^{-4}$ | $1.19 \times 10^{-4}$ | $-10.2\%$ | 0.16% |
| $\tilde{\chi}_1^+ \rightarrow \tau^+ \tilde{\nu}_\tau$ | $1.47 \times 10^{-2}$ | $1.48 \times 10^{-2}$ | $+0.1\%$ | 19.58% |
| $\tilde{\chi}_1^+ \rightarrow \mu^+ \tilde{\nu}_\mu$ | $1.06 \times 10^{-2}$ | $1.07 \times 10^{-2}$ | $+1.0\%$ | 14.24% |
| $\tilde{\chi}_1^+ \rightarrow e^+ \tilde{\nu}_e$ | $1.06 \times 10^{-2}$ | $1.07 \times 10^{-2}$ | $+1.0\%$ | 14.22% |
| $\tilde{\chi}_1^+ \rightarrow W^+\tilde{\chi}_1^0$ | $9.65 \times 10^{-4}$ | $1.28 \times 10^{-3}$ | $+32.3\%$ | 1.69% |

Table 5: One-loop corrections on $\tilde{\chi}_1^+$ decay widths for set (A)

In Table 7, we show the one-loop electroweak and QCD corrections to the various 3-body decay widths of $\tilde{\chi}_1^+$ in the parameter set (B), for which the two-body decays of $\tilde{\chi}_1^+$ are kinematically forbidden. The three-body decay widths of $\tilde{\chi}_2^+$ are not shown, since, decaying dominantly into two bodies, it has extremely small three-body decay branching ratios. Note that for the decay modes involving quarks $\tilde{\chi}_1^+ \rightarrow q\bar{q}\tilde{\chi}_1^0$, the electroweak and the QCD corrections are separately given.

We should note that the lighter chargino $\tilde{\chi}_1^+$ is the next lightest SUSY-particle (NLSP) in the set (B). Since NLSP must be first produced by acceler-
Table 6: One-loop corrections on $\tilde{\chi}_2^+$ decay widths for set (A)

| $\tilde{\chi}_2^+$ decay | $\Gamma_0$ (GeV) | $\Gamma$ (GeV) | $\delta \Gamma/\Gamma_0$ | Br |
|--------------------------|------------------|----------------|------------------------|----|
| $\tilde{\chi}_2^+ \rightarrow \nu_e \tilde{\tau}_2^+$ | $1.54 \times 10^{-1}$ | $1.48 \times 10^{-1}$ | $-3.9\%$ | $4.20\%$ |
| $\tilde{\chi}_2^+ \rightarrow \nu_\mu \tilde{\mu}_2^+$ | $1.36 \times 10^{-1}$ | $1.46 \times 10^{-1}$ | $+7.5\%$ | $4.13\%$ |
| $\tilde{\chi}_2^+ \rightarrow \nu_e \tilde{e}_2^+$ | $1.36 \times 10^{-1}$ | $1.46 \times 10^{-1}$ | $+7.6\%$ | $4.14\%$ |
| $\tilde{\chi}_2^+ \rightarrow \tau^+ \tilde{\nu}_\tau$ | $6.89 \times 10^{-2}$ | $5.70 \times 10^{-2}$ | $-17.3\%$ | $1.61\%$ |
| $\tilde{\chi}_2^+ \rightarrow \mu^+ \tilde{\nu}_\mu$ | $4.33 \times 10^{-2}$ | $5.38 \times 10^{-2}$ | $+24.2\%$ | $1.52\%$ |
| $\tilde{\chi}_2^+ \rightarrow e^+ \tilde{\nu}_e$ | $4.32 \times 10^{-2}$ | $5.37 \times 10^{-2}$ | $+24.4\%$ | $1.52\%$ |
| $\tilde{\chi}_2^+ \rightarrow W^+ \tilde{\chi}_1^0$ | $1.93 \times 10^{-1}$ | $2.07 \times 10^{-1}$ | $+7.0\%$ | $5.85\%$ |
| $\tilde{\chi}_2^+ \rightarrow W^+ \tilde{\chi}_2^0$ | $8.66 \times 10^{-1}$ | $9.93 \times 10^{-1}$ | $+14.6\%$ | $28.12\%$ |
| $\tilde{\chi}_2^+ \rightarrow Z \tilde{\chi}_1^0$ | $7.53 \times 10^{-1}$ | $8.56 \times 10^{-1}$ | $+13.7\%$ | $24.26\%$ |
| $\tilde{\chi}_2^+ \rightarrow h^0 \tilde{\chi}_1^0$ | $5.97 \times 10^{-1}$ | $6.07 \times 10^{-1}$ | $+1.75\%$ | $17.20\%$ |
| $\tilde{\chi}_2^+ \rightarrow \bar{b} \tilde{t}_1$ | $2.82 \times 10^{-1}$ | $2.57 \times 10^{-1}$ | $\left\{ \begin{array}{l} -8.9\%(\text{ELWK}) \\ +1.8\%(\text{QCD}) \end{array} \right.$ | $7.43\%$ |

By combining the production cross section (Fig.1) and the decay branching ratios (Table 7) for the set (B), we obtain the one-loop corrected cross sections for the direct experimental signals. Fig.2 shows the energy dependence of the cross section for the two types of the chargino signals, $e^+e^-$ + missing energies and 4-jets + missing energies, at $e^+e^-$ colliders for the set (B).
Table 7: One-loop corrections on $\tilde{\chi}_1^+$ decay widths for set (B)

| Decay | $\Gamma_0$ (GeV) | $\Gamma$ (GeV) | $\delta \Gamma / \Gamma_0$ | $\text{Br}$ |
|-------|-----------------|----------------|---------------------|---------|
| $\tilde{\chi}_1^+ \rightarrow e^+ \nu \tilde{\chi}_1^0$ | $4.42 \times 10^{-6}$ | $4.48 \times 10^{-6}$ | +9.4% | 20.18% |
| $\tilde{\chi}_1^+ \rightarrow \mu^+ \nu \tilde{\chi}_1^0$ | $4.42 \times 10^{-6}$ | $4.48 \times 10^{-6}$ | +9.4% | 20.18% |
| $\tilde{\chi}_1^+ \rightarrow \tau^+ \nu \tilde{\chi}_1^0$ | $6.46 \times 10^{-6}$ | $7.22 \times 10^{-6}$ | +11.8% | 30.09% |
| $\tilde{\chi}_1^+ \rightarrow u \bar{d} \tilde{\chi}_1^0$ | $3.35 \times 10^{-6}$ | $3.55 \times 10^{-6}$ | $-0.2\%$(ELWK) +6.3%(QCD) | 14.81% |
| $\tilde{\chi}_1^+ \rightarrow c \bar{s} \tilde{\chi}_1^0$ | $3.33 \times 10^{-6}$ | $3.54 \times 10^{-6}$ | $-0.2\%$(ELWK) +6.3%(QCD) | 14.74% |

5 Comments

The total cross section is a sum of the one-loop corrected cross section and the cross section of the hard photon radiation. The former is negative and large, while the latter is positive and large, originating mainly from the initial radiation. In our calculation for the hard photon radiation, we do not set an energy cut in the upper limit nor any angle cut.

In the calculation of the one-loop correction of the decay widths the situation is almost the same except that in the evaluation of the decay widths the contribution of the real photon (gluon) emission from both initial and final states is important. The correction proportional to the fermion mass $m_f$ is expected to be large for the third generation $\tau$, $t$ and $b$, and an additional enhancement can emerge in the case of large $\tan \beta$ for $\tau$ and $b$. Note that we set $\tan \beta = 10$ for both sets (A) and (B). We see these effects in the numerical results presented in Tables 5, 6 and 7, where the branching ratios of the $\tau$ modes are larger than $e$ and $\mu$ modes.

The radiative correction on the chargino pair production in the SPA1a' scenario (set (A)) has been studied by the Wien group [27]. Unfortunately, as they adopt different treatment of the photon emission correction, the direct comparison of our result for the full electroweak correction with theirs is not possible. For the comparison we extract the "weak" correction $\Delta \sigma_{\text{weak}}$ defined by

$$\Delta \sigma_{\text{weak}} \equiv \sigma_{\text{elwk}} - \sigma_{\text{BORN}} * \delta_{\text{QED}} - \sigma_{\text{hard}}^\text{initial}. \quad (5.1)$$

We find that the result of GRACE/SUSY-loop is consistent with the previous
Figure 1: Total cross section for chargino pair production $e^+e^- \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$ for the set (A) and set (B). Solid line and dotted line denotes $\sigma_{1\text{loop}}$ and $\sigma_{\text{BORN}}$, respectively.

In the parameter setting (A), the produced chargino dominantly decays into two bodies, because the mass of the chargino $m_{\tilde{\chi}_1^+} (= 184.2\text{GeV})$ is larger than the mass of the lighter charged sleptons, $m_{\tilde{\ell}_1}$, and $m_{\tilde{\nu}_1}$ as well as $m_W + m_{\tilde{\chi}_1^0}$ (see Tables 1 and 3). Since the chargino cannot decay into any squark in two-body decay modes and $\text{BR}(\tilde{\chi}_1^+ \rightarrow W^+\tilde{\chi}_1^0 \rightarrow q\bar{q}'\tilde{\chi}_1^0)$ is negligibly small (see Table 5), we cannot expect the signals with the quark-jets from the $\tilde{\chi}_1^+$ pair production. We find also that $\text{BR}(\tilde{\ell}_1 \rightarrow \ell\tilde{\chi}_1^0) = 1$ and $\text{BR}(\tilde{\nu}_1 \rightarrow \nu\tilde{\chi}_1^0) = 1$ in the parameter setting (A). This means that the most plausible experimental signal of the chargino pair production is the lepton pair plus the large missing energies. The precise measurement of the energy and the momentum of the $\tau$ leptons is particularly important because $\tau^+\tau^-$ signal is the dominant mode in this case.
Figure 2: Total cross section for $e^+e^- \rightarrow \tilde{\chi}_1^+\tilde{\chi}_1^- \rightarrow (e^+\nu_e\tilde{\chi}_0^0)(e^-\bar{\nu}_e\tilde{\chi}_0^0)$ (a) and $e^+e^- \rightarrow \tilde{\chi}_1^+\tilde{\chi}_1^- \rightarrow (q\bar{q}'\tilde{\chi}_0^0)(q\bar{q}'\tilde{\chi}_0^0)$ (b). Solid line and dotted line denotes $\sigma_{\text{1-loop}}$ and $\sigma_{\text{BORN}}$, respectively.

In the parameter setting (B), the chargino $\tilde{\chi}_1^+$ with $m_{\tilde{\chi}_1^+} = 147.0\text{GeV}$ is lighter than in set (A) in which $m_{\tilde{\chi}_1^+} = 184.2\text{GeV}$. We find the apparent shift of the production peak to the lower energy due to the different chargino mass value used in the set (A) and (B). As has been discussed in the previous section, two-body decays of $\tilde{\chi}_1^+$ are kinematically forbidden in the set (B). We can use both types of signals, the lepton pair plus large missing energies and the quark-jets plus large missing energies, for the chargino detection.

We have developed a tool for the full automatic one-loop calculation of the MSSM processes, GRACE/SUSY-1loop, which is characterized by the gauge symmetric and the on-shell renormalization scheme, and includes the various self-consistency check schemes. It is certainly a useful tool for the present and future precise analyses of the experimental data.

Acknowledgement
We would like to thank Y. Shimizu, T. Kaneko, G. Bélanger and F. Boudjema for fruitful discussions. This work was partly supported by Japan Society
for the Promotion of Science under Grant-in-Aid for Scientific Research (B) (No.17340085).
Appendix A. Expressions of counterterms

In this Appendix, we list the expression of the counterterms and the renormalization constants in terms of two-point functions.

The wavefunction renormalization constants are expanded in the one-loop order as
\[
Z^{1/2}_X = 1 + \frac{1}{2} \delta Z_X, \quad (A.1)
\]
\[
Z^{1/2}_{XY} = \begin{cases} 
1 + \frac{1}{2} \delta Z_{XY}, & X = Y \\
\frac{1}{2} \delta Z_{XY}, & X \neq Y
\end{cases} \quad (A.2)
\]

We use the following abbreviations;
\[
c_W = \cos \theta_W, \quad c_L = \cos \phi_L, \quad c_R = \cos \phi_R, \quad c_\alpha = \cos \alpha, \quad c_\beta = \cos \beta,
\]
\[
s_W = \sin \theta_W, \quad s_L = \sin \phi_L, \quad s_R = \sin \phi_R, \quad s_\alpha = \sin \alpha, \quad s_\beta = \sin \beta \quad (A.3)
\]

where \(\phi_L\) and \(\phi_R\) are the mixing angles which diagonalize the chargino mass matrix (see (2.10) of [11]), and \(\alpha\) is the mixing angle in the CP even Higgs sector.

Other counterterms appearing in the gauge sector are expressed in terms of \(\delta Z_W, \delta Z_B\text{ and }\delta Z_g\). For example,
\[
\frac{\delta g}{g} = \delta Z_g - \frac{3}{2} \delta Z_W, \quad (A.7)
\]
\[
\frac{\delta g'}{g'} = \frac{1}{2} \delta Z_B, \quad (A.8)
\]
\[
\frac{\delta s_W}{s_W} = -c_W^2 (\frac{\delta g}{g} - \frac{\delta g'}{g'}), \quad (A.9)
\]
\[
\frac{\delta c_W}{c_W} = +s_W^2 (\frac{\delta g}{g} - \frac{\delta g'}{g'}). \quad (A.10)
\]
fermion sector

\[ \delta m_f = -Re\Sigma_f^S(m_f) - m_f Re\Sigma_f^V(m_f), \quad (A.11) \]
\[ \delta Z_f^R = Re\Sigma_f^V(m_f) + Re\Sigma_f^A(m_f) \]
\[ + 2m_f [ Re\Sigma_f^S(m_f) + m_f Re\Sigma_f^V(m_f) ], \quad (A.12) \]
\[ \delta Z_f^L = Re\Sigma_f^V(m_f) - Re\Sigma_f^A(m_f) \]
\[ + 2m_f [ Re\Sigma_f^S(m_f) + m_f Re\Sigma_f^V(m_f) ]. \quad (A.13) \]

where the selfenergy function of Dirac fermion \( f \) is decomposed as

\[ \Sigma_f(q^2) \equiv \Sigma_f^S(q^2) 1 + \Sigma_f^P(q^2) \gamma_5 + \Sigma_f^V(q^2) \gamma_\mu + \Sigma_f^A(q^2) \gamma_\mu \gamma_5. \quad (A.14) \]

tadpole terms

\[ \delta T_i = T_{\phi_i}^{\text{loop}}, \quad i = 1, 2, \quad (A.15) \]

where \( \delta T_i \) is the tadpole counterterms.

Higgs sector

\[ \delta M_A^2 = -Re\Sigma_{A^0A^0}(M_A^2) + M_A^2 Re\Sigma_{A^0A^0}'(M_A^2), \quad (A.16) \]
\[ \left( \begin{array}{c}
\delta Z_{H_1} \\
\delta Z_{H_2} \\
\delta \nu_1 \\
\delta \nu_2 \\
v_1 \\
v_2 
\end{array} \right) = 
\left( \begin{array}{ccc}
2c_\beta^2 & -2s_\beta^2 & 0 \\
-2s_\beta^2 & c_\beta^2 & 0 \\
c_\beta & -s_\beta & C(i, j) + \frac{1}{2} \\
s_\beta & c_\beta & 0
\end{array} \right) 
\left( \begin{array}{c}
\frac{1}{c_\beta s_\beta} Re\Sigma_{A^0A^0}'(M_A^2) \\
C\delta X \\
C_2\delta \phi Y
\end{array} \right), \quad (A.17) \]

where

\[ C = \frac{2}{\sin 2\beta (\sin 2\beta + \sin 2\alpha)}, \]
\[ C(1, 2) = 3s_\alpha^2c_\beta^2 - s_\alpha^2 c_\beta^2 + 2s_\alpha c_\alpha s_\beta c_\beta, \]
\[ C(1, 3) = c_\beta(c_\alpha^2 - s_\beta^2), \]
\[ C(2, 2) = 3s_\beta^2c_\alpha^2 - c_\beta^2 s_\alpha^2 + 2s_\alpha c_\alpha s_\beta c_\beta, \]
\[ C(2, 3) = -s_\beta^2(c_\alpha^2 - s_\beta^2), \]
\[ C(3, 2) = \frac{1}{2} [c_\beta^2(2s_\beta^2 - 1) + c_\alpha^2 + 2s_\alpha c_\alpha s_\beta c_\beta], \]
\[ C(3,3) = -\frac{1}{2}[c_\beta(2s_\beta^2 - 1) + s_\alpha^2 + 2s_\alpha c_\alpha s_\beta c_\beta], \]
\[ C(4,2) = +\frac{1}{2}[s_\beta(2c_\beta^2 - 1) + s_\alpha^2 + 2s_\alpha c_\alpha s_\beta]. \]
\[ C(4,3) = -\frac{1}{2}[s_\beta(2c_\beta^2 - 1) + c_\alpha^2 + 2s_\alpha c_\alpha s_\beta]. \] (A.18)

and

\[
\delta X = \text{Re}[\Pi'_{AA}(0) + 2c_W s_W \Pi_{AZ}(0) M_W^2 + (s_W^2 s_W^2 \Pi W(M_W^2) - c_W^4 \Pi ZZ(M_Z^2)] M^2 W, \]
\[
\delta Y = -\frac{1}{M_{H^0}^2} \left[ \text{Re}[\Sigma H^0 H^0(M_{H^0}^2) + \sin^2(\alpha - \beta)\delta M_A^2 + \cos^2(\alpha + \beta)M_Z^2 \delta Z Z \right.
\]
\[
+ \frac{g}{2M_W} \cos(\alpha - \beta)[1 + \sin^2(\alpha - \beta)]\delta T_{H^0}^0
\]
\[
+ \frac{g}{2M_W} \sin(\alpha - \beta) \cos^2(\alpha - \beta)\delta T_{H^0}^0 \right], \] (A.19)

with

\[ \delta Z_Z = 2c_\beta^2 \delta Z_g - 3c_\beta^2 \delta Z_W - s_W^2 \delta Z_B. \] (A.21)

The tadpole counterterms for the physical Higgs bosons are defined by

\[
\begin{pmatrix}
\delta T_{H^0}^0 \\
\delta T_{H^0}^1
\end{pmatrix} = \begin{pmatrix}
\cos \alpha & \sin \alpha \\
-\sin \alpha & \cos \alpha
\end{pmatrix}
\begin{pmatrix}
\delta T_1 \\
\delta T_2
\end{pmatrix}. \] (A.22)

The counterterm of \( \tan \beta \) is then determined by

\[
\delta \tan \beta = -\frac{1}{2} \tan \beta(\delta Z_{H_1} - \delta Z_{H_2} - 2\frac{\delta v_1}{v_1} + 2\frac{\delta v_2}{v_2}), \] (A.23)

while the gauge-boson mass counterterms are given by

\[
\delta M_W^2 = M_W^2 [2\delta Z_g - 3\delta Z_W + \cos^2 \beta \delta Z_{H_1} + \sin^2 \beta \delta Z_{H_2} - 2 \cos^2 \beta \frac{\delta v_1}{v_1} - 2 \sin^2 \beta \frac{\delta v_2}{v_2}], \] (A.24)
\[
\delta M_Z^2 = M_Z^2 [\delta Z_Z + \cos^2 \beta \delta Z_{H_1} + \sin^2 \beta \delta Z_{H_2} - 2 \cos^2 \beta \frac{\delta v_1}{v_1} - 2 \sin^2 \beta \frac{\delta v_2}{v_2}]. \] (A.25)
Chargino sector

\[
\begin{pmatrix}
\delta Z_{\chi^w} \\
\delta Z_{\tilde{H}_1} \\
\delta Z_{\tilde{H}_2}
\end{pmatrix}
= \begin{pmatrix}
1 & (s^2_R c_R + c^2_L s^2_R)/X & -2 s^2_L s^2_R/X \\
1 & (2s^2_L c^2_R - c^2_L - c^2_R)/X & 2s^2_L s^2_R/X \\
1 & (2c^2_L s^2_R - c^2_L - c^2_R)/X & 2c^2_L s^2_R/X
\end{pmatrix}
\begin{pmatrix}
A \\
Re\Sigma_{11}^A(m_{\tilde{\chi}^1}) \\
Re\Sigma_{22}^A(m_{\tilde{\chi}^2})
\end{pmatrix},
\]

\( (A.26) \)

\[
\delta M_2 = \frac{s_L s_R(\Delta_{22} - \epsilon_L\Pi_{22}) - c_L c_R(\Delta_{11} - \Pi_{11})}{c^2_L - s^2_R}, \quad (A.27)
\]

\[
\delta \mu = \frac{s_L s_R(\Delta_{11} - \Pi_{11}) - c_L c_R(\Delta_{22} - \epsilon_L\Pi_{22})}{c^2_L - s^2_R}, \quad (A.28)
\]

where

\[
X = s^2_L - s^2_R, \quad (A.29)
\]

\[
A = 2m_{\tilde{\chi}^1} [Re\Sigma_{11}^S(m_{\tilde{\chi}^1}) + m_{\tilde{\chi}^1} Re\Sigma_{11}^V(m_{\tilde{\chi}^1})] + Re\Sigma_{11}^V(m_{\tilde{\chi}^1}), \quad (A.30)
\]

\[
\Delta_{11} = + \left( c_L c_R M_2 + c_L s_R M_W \frac{\cos \beta}{\sqrt{2}} + s_L c_R M_W \frac{\sin \beta}{\sqrt{2}} \right) \delta Z_{\chi^w}
\]

\[
+ \left( c_L s_R \sqrt{2} M_W \cos \beta + s_L s_R \mu \right) \frac{1}{2} \delta Z_{\tilde{H}_1}
\]

\[
+ \left( s_L c_R \sqrt{2} M_W \sin \beta + s_L s_R \mu \right) \frac{1}{2} \delta Z_{\tilde{H}_2}
\]

\[
+ c_L s_R \sqrt{2} M_W \cos \beta \left( \frac{1}{2} \delta Z_{H_1} + \frac{\delta g}{g} - \frac{\delta v_1}{v_1} \right)
\]

\[
+ s_L c_R \sqrt{2} M_W \sin \beta \left( \frac{1}{2} \delta Z_{H_2} + \frac{\delta g}{g} - \frac{\delta v_2}{v_2} \right), \quad (A.31)
\]

\[
\Delta_{22} = + \left( s_L s_R M_2 - s_L c_R M_W \frac{\cos \beta}{\sqrt{2}} - c_L s_R M_W \frac{\sin \beta}{\sqrt{2}} \right) \delta Z_{\chi^w}
\]

\[
+ \left( -s_L c_R \sqrt{2} M_W \cos \beta + c_L c_R \mu \right) \frac{1}{2} \delta Z_{\tilde{H}_1}
\]

\[
+ \left( -c_L s_R \sqrt{2} M_W \sin \beta + c_L c_R \mu \right) \frac{1}{2} \delta Z_{\tilde{H}_2}
\]

\[
- s_L c_R \sqrt{2} M_W \cos \beta \left( \frac{1}{2} \delta Z_{H_1} + \frac{\delta g}{g} - \frac{\delta v_1}{v_1} \right)
\]

\[
- c_L s_R \sqrt{2} M_W \sin \beta \left( \frac{1}{2} \delta Z_{H_2} + \frac{\delta g}{g} - \frac{\delta v_2}{v_2} \right), \quad (A.32)
\]

\[
\Pi_{11} = 2m^2_{\chi^1} [Re\Sigma_{11}^S(m_{\chi^1}) + m_{\chi^1} Re\Sigma_{11}^V(m_{\chi^1})] - Re\Sigma_{11}^S(m_{\chi^1}), \quad (A.33)
\]
\[ \Pi_{22} = \frac{m_{\tilde{\chi}^2}}{2} \left[ (s_l^2 + s_R^2) \delta Z_{\lambda^w} + c_R^2 \delta Z_{H_1} + c_L^2 \delta Z_{H_2} \right] - \text{Re} \Sigma_2^S (m_{\tilde{\chi}^2}^+) - m_{\tilde{\chi}^2}^+ \text{Re} \Sigma_2^V (m_{\tilde{\chi}^2}^+). \]  

(A.34)

Neutralino sector

\[ \delta Z_{\lambda} = \frac{1}{\mathcal{O}_{11}^2} \left[ \delta Z_{11} - \mathcal{O}_{12}^2 \delta Z_{\lambda^w} - \mathcal{O}_{13}^2 \delta Z_{H_1} - \mathcal{O}_{14}^2 \delta Z_{H_2} \right] \]  

(A.35)

\[ \delta M_1 = \frac{1}{\mathcal{O}_{11}^2} \left[ \delta m_{11} - \sum_{(p,q) \neq (1,1)} \mathcal{O}_{1p} \mathcal{O}_{1q} (\delta M_N)_{pq} - \sum_{(p,q)} \mathcal{O}_{1p} \mathcal{O}_{1q} \delta Z_p (M_N)_{pq} \right] \]  

(A.36)

where \( \mathcal{O} \) is the orthogonal matrix which diagonalizes the neutralino mass matrix \( M_N \)

\[ \mathcal{O} M_N \mathcal{O}^t = \begin{pmatrix} m_{\tilde{\chi}_1} & \phantom{-}m_{\tilde{\chi}_2} & m_{\tilde{\chi}_3} & m_{\tilde{\chi}_4} \\ \end{pmatrix} \]  

(A.37)

and

\[ \delta M_N = \begin{pmatrix} \delta M_1 & 0 & -M_Z s_W \cos \beta \Delta_{13} & M_Z s_W \sin \beta \Delta_{14} \\
* & \delta M_2 & M_Z c_W \cos \beta \Delta_{23} & -M_Z c_W \sin \beta \Delta_{24} \\
* & * & 0 & -\delta \mu \\
* & * & * & 0 \end{pmatrix}, \]  

(A.38)

with

\[ \Delta_{13} = \frac{\delta M_Z^2}{2M_Z^2} - c_Z^2 \left( \frac{\delta g}{g} - \frac{\delta g'}{g'} \right) - \tan \beta \cos^2 \beta \delta \tan \beta, \]  

(A.39)

\[ \Delta_{14} = \frac{\delta M_Z^2}{2M_Z^2} - s_Z^2 \left( \frac{\delta g}{g} - \frac{\delta g'}{g'} \right) + \cot \beta \cos^2 \beta \delta \tan \beta, \]  

(A.40)

\[ \Delta_{23} = \frac{\delta M_Z^2}{2M_Z^2} + s_Z^2 \left( \frac{\delta g}{g} - \frac{\delta g'}{g'} \right) - \tan \beta \cos^2 \beta \delta \tan \beta, \]  

(A.41)

\[ \Delta_{24} = \frac{\delta M_Z^2}{2M_Z^2} + s_Z^2 \left( \frac{\delta g}{g} - \frac{\delta g'}{g'} \right) + \cot \beta \cos^2 \beta \delta \tan \beta. \]  

(A.42)

The phase factor \( \eta_i \) in (A.36) is to convert the negative mass eigenvalue in (A.37) to positive,

\[ \eta_i = \begin{cases} 1 & m_{\tilde{\chi}_i} > 0 \\
i & m_{\tilde{\chi}_i} < 0 \end{cases}. \]  

(A.43)
sfermion sector
For simplicity, we show only the expression for the first generation.

\[
\delta m_f^2 = -\text{Re} \Sigma_{ff}(m_f^2), \quad f = u_1, u_2, d_1, d_2, e_1, e_2, \nu_e, \tag{A.44}
\]

\[
\delta Z_{ff} = \Sigma'_{ff}(m_f^2), \quad f = u_1, u_2, d_1, d_2, e_1, e_2, \nu_e, \tag{A.45}
\]

\[
\frac{1}{2} \delta Z_{fi,fj} = -\frac{\Sigma'_{fi,fj}(m_f^2)}{m_i^2 - m_j^2}, \quad i \neq j, \quad f = u, d, e. \tag{A.46}
\]

\[
\delta \theta_e = \frac{\delta m_{\nu_e} - \delta (M_W^2 \cos 2\beta - m_e^2) - \cos^2 \theta_e \delta m_{\tilde{e}_1}^2 - \sin^2 \theta_e \delta m_{\tilde{e}_2}^2}{\sin 2\theta_e (m_{\tilde{e}_2}^2 - m_{\tilde{e}_1}^2)}, \tag{A.47}
\]

\[
\delta \theta_u = \frac{1}{2} \Sigma_{\tilde{u}_1, \tilde{u}_2}(m_{\tilde{u}_1}^2) + \Sigma_{\tilde{u}_1, \tilde{u}_2}(m_{\tilde{u}_1}^2)}{m_{\tilde{u}_2}^2 - m_{\tilde{u}_1}^2} \tag{A.48}
\]

\[
\delta \theta_d = \frac{\delta (\cos^2 \theta_e m_{\tilde{u}_1}^2 + \sin^2 \theta_e m_{\tilde{u}_2}^2 - M_W^2 \cos 2\beta - m_u^2 + m_d^2) - \cos^2 \theta_d \delta m_{d_1}^2 - \sin^2 \theta_d \delta m_{d_2}^2}{\sin 2\theta_d (m_{d_2}^2 - m_{d_1}^2)}, \tag{A.49}
\]

QCD sector

\[
\delta M_3 = -m_\tilde{g}[\text{Re}\Sigma_{\tilde{g}}(m_\tilde{g}) + \text{Re}\Sigma'_{\tilde{g}}(m_\tilde{g})], \tag{A.50}
\]

\[
\delta Z_{\tilde{g}} = \text{Re}\Sigma'_{\tilde{g}}(m_\tilde{g}) + 2m_\tilde{g}^2[\text{Re}\Sigma_{\tilde{g}}(m_\tilde{g}) + \text{Re}\Sigma'_{\tilde{g}}(m_\tilde{g})], \tag{A.51}
\]

\[
\delta Z_{\text{gluon}} = \text{Re}\Pi'_{gg}(0), \tag{A.52}
\]

\[
\delta Z_{g_s} = -C\left[\frac{2}{4 - d} - \gamma_E + \ln(4\pi)\right]. \tag{A.53}
\]

where \(C\) is the finite constant appearing at the one-loop vertex correction as

\[
(iV_{\mu\nu\lambda})[C(\frac{2}{4 - d} - \gamma_E + \ln(4\pi)) + \ldots]. \tag{A.54}
\]
Appendix B. External wavefunction renormalization constants

We list the external wavefunction renormalization constant $\delta Z^{ext}$ which appears in the amplitude as

$$\mathcal{M} \sim \frac{1}{2} \delta Z^{ext} \times \text{(Born amplitude)} \quad (B.1)$$

**Gauge bosons**

$$\delta Z^{ext}_W = \hat{\Pi}'_W(M^2_W) = \Pi'_W(M^2_W) - \delta Z_W, \quad (B.2)$$

$$\delta Z^{ext}_Z = \hat{\Pi}'_{ZZ}(M^2_Z) = \Pi'_{ZZ}(M^2_Z) - \delta Z_{ZZ}, \quad (B.3)$$

**Higgs sector**

$$\delta Z^{ext}_{h^0} = \hat{\Sigma}'_{h^0}(M^2_{h^0}). \quad (B.4)$$

Since the pole mass of $h^0$ and $H^\pm$ in one-loop order does not agree with the tree mass, some complication appears.

$$\delta Z^{ext}_{h^0} = \frac{1}{1 - \hat{\Sigma}'_{hh}(M^2_{h^0}, m^2_{h^0}, m^2_{H^0})} - 1 \approx \hat{\Sigma}'_{hh}(M^2_{h^0}, m^2_{h^0}, m^2_{H^0}) \quad (B.5)$$

where

$$\hat{\Sigma}'_{hh}(M^2_{h^0}) = \frac{\partial}{\partial q^2} \hat{\Sigma}_{hh}(q^2)\big|_{M^2_{h^0}} + \frac{\partial}{\partial q^2} \left[ q^2 - m^2_{H^0} - \hat{\Sigma}_{HH}(q^2) \right]\big|_{M^2_{h^0}}. \quad (B.6)$$

and $M_{h^0}$ is the one-loop improved pole mass of $h^0$. The renormalized self-energy functions are defined by

$$\hat{\Sigma}(q^2)_{hh} = \Sigma(q^2)_{hh} + \delta M^2_{hh} - q^2 (\sin^2 \alpha \Re \delta Z_{H_1} + \cos^2 \alpha \Re \delta Z_{H_2}), \quad (B.7)$$

$$\hat{\Sigma}(q^2)_{HH} = \Sigma(q^2)_{HH} + \delta M^2_{HH} - q^2 (\cos^2 \alpha \Re \delta Z_{H_1} + \sin^2 \alpha \Re \delta Z_{H_2}), \quad (B.8)$$

$$\hat{\Sigma}(q^2)_{Hh} = \Sigma(q^2)_{Hh} + \delta M^2_{Hh} - q^2 \cos \alpha \sin \alpha (\Re \delta Z_{H_2} - \Re \delta Z_{H_1}), \quad (B.9)$$
with
\[
\delta M_{hh}^2 = \cos^2(\alpha - \beta)\delta M_A^2 - \frac{g}{2M_W} \cos(\alpha - \beta)\sin^2(\alpha - \beta)T_{h_0}^{\text{loop}}
\]
\[
- \frac{g}{2M_W} \sin(\alpha - \beta)[1 + \cos^2(\alpha - \beta)]T_{h_0}^{\text{loop}} + M_E^2[\sin^2(\alpha + \beta)(\delta Z_Z + \delta Z_{H_1} + \delta Z_{H_2})
\]
\[
+ \sin(\alpha + \beta)\sin(\alpha - \beta)(\delta Z_{H_1} - \delta Z_{H_2})] + 2\sin(\alpha + \beta)[\sin \beta \cos \beta \cos(\alpha + \beta) - \sin \alpha \cos \beta]M_Z^2 \frac{\delta v_1}{v_1}
\]
\[
-2\sin(\alpha + \beta)[\sin \beta \cos \beta \cos(\alpha + \beta) + \cos \alpha \sin \beta]M_Z^2 \frac{\delta v_2}{v_2},
\]
\[
\delta M_{HH}^2 = -Re\Sigma_{HH}(M_H^2) + M_{H_0}^2(\cos^2 \alpha \delta Z_{H_1} + \sin^2 \alpha \delta Z_{H_2}),
\]
\[
\delta M_{Hh}^2 = -\sin(\beta - \alpha)\cos(\beta - \alpha)\delta M_A^2 + \frac{g}{2M_W} \sin^3(\beta - \alpha)T_{h_0}^{\text{loop}}
\]
\[
+ \frac{g}{2M_W} \cos^3(\alpha - \beta)T_{h_0}^{\text{loop}} - M_E^2[\sin(\alpha + \beta)\cos(\alpha + \beta)(\delta Z_Z + \delta Z_{H_1} + \delta Z_{H_2})
\]
\[
+ \sin \alpha \cos \alpha(\delta Z_{H_1} - \delta Z_{H_2})] + \frac{1}{2} \sin(2\alpha + 2\beta)(1 + \cos 2\beta) - \frac{\sin 2\alpha \cos 2\alpha \sin 2\beta}{\sin(2\alpha - 2\beta)} M_Z^2 \frac{\delta v_1}{v_1}
\]
\[
+ \frac{1}{2} \sin(2\alpha + 2\beta)(1 - \cos 2\beta) + \frac{\sin 2\alpha \cos 2\alpha \sin 2\beta}{\sin(2\alpha - 2\beta)} M_Z^2 \frac{\delta v_2}{v_2}.
\]

The expressions \([B.5]\) and \([B.6]\) agree with those given in \([26]\).

\[
\delta Z_{H \pm}^{\text{ext}} = \frac{1}{1 - \hat{\Sigma}_{H \pm H \pm}^{\text{st}}(M_{H \pm}^2)} - 1 \sim \hat{\Sigma}_{H \pm H \pm}^{\text{st}}(M_{H \pm}^2),
\]

where
\[
\hat{\Sigma}_{H \pm H \pm}^{\text{st}}(q^2) = \frac{\partial}{\partial q^2} \hat{\Sigma}_{H \pm H \pm}(q^2)\bigg|_{M_{H \pm}^2}
\]
\[
+ \frac{\partial}{\partial q^2} \left[ \frac{\hat{\Sigma}_{H \pm G \pm}(q^2)}{q^2 - m_{H \pm}^2 - \hat{\Sigma}_{H \pm H \pm}(q^2)} \right] \bigg|_{M_{H \pm}^2}.
\]

The renormalized selfenergy functions appearing in \([B.14]\) are given by
\[
\hat{\Sigma}(q^2)_{H \pm H \pm} = \Sigma(q^2)_{H \pm H \pm} + \delta M_{H \pm H \pm}^2 - q^2(\delta Z_{H_1} + \delta Z_{H_2}), \quad \Sigma(q^2)_{H \pm G \pm} = \Sigma(q^2)_{H \pm G \pm} + \delta M_{H \pm G \pm}^2 - q^2(\delta Z_{H_2} - \delta Z_{H_1}),
\]

25
with
\[
\delta M^2_{H^\pm H^\pm} = \delta M^2_{A^0 A^0} + M^2_W (\delta Z_x - 2 c_\beta \frac{\delta v_1}{v_1} - 2 s_\beta \frac{\delta v_2}{v_2}), \quad (B.17)
\]
\[
\delta M^2_{H^\pm G^\pm} = \delta M^2_{G^0 A^0} - c_\beta s_\beta M^2_W (\frac{\delta v_1}{v_1} - \frac{\delta v_2}{v_2}). \quad (B.18)
\]

In the one-loop order, (B.5) and (B.14) become
\[
\delta Z^{\text{ext}}_{hh} = \hat{\Sigma}^{\prime}_{hh} (M^2_{h^0}), \quad (B.19)
\]
\[
\delta Z^{\text{ext}}_{H^\pm H^\pm} = \hat{\Sigma}^{\prime}_{H^\pm H^\pm} (M^2_{H^\pm}). \quad (B.20)
\]

**Chargino**
\[
\delta Z^{\text{ext}}_{\tilde{\chi}^+_{1}} = 2 M^2_{\tilde{\chi}^+_{1}} (\Sigma^S_{\tilde{\chi}^+_{1}} (M^2_{\tilde{\chi}^+_{1}}) + \Sigma^V_{\tilde{\chi}^+_{1}} (M^2_{\tilde{\chi}^+_{1}})) + \Sigma^V_{\tilde{\chi}^+_{1}} (M^2_{\tilde{\chi}^+_{1}}) - \frac{1}{2} (\delta Z^R_{22} + \delta Z^L_{22}), \quad (B.21)
\]
where the chargino selfenergy functions are decomposed as (A.14). In terms of the renormalization constants introduced in section 2, the chargino wavefunction renormalization constants are given by
\[
\delta Z^L_{22} = \sin \phi^2 L \delta Z_{\lambda^0} + \cos \phi^2 L \delta Z_{\tilde{H}_2},
\]
\[
\delta Z^R_{22} = \sin \phi^2 R \delta Z_{\lambda^0} + \cos \phi^2 R \delta Z_{\tilde{H}_1}. \quad (B.22)
\]

**Neutralino**
\[
\delta Z^{\text{ext}}_{\tilde{\chi}^0_i} = 2 M^2_{\tilde{\chi}^0_i} [\Sigma^S_{\tilde{\chi}^0_i} (M^2_{\tilde{\chi}^0_i}) + \Sigma^V_{\tilde{\chi}^0_i} (M^2_{\tilde{\chi}^0_i})] + \Sigma^V_{\tilde{\chi}^0_i} (M^2_{\tilde{\chi}^0_i}) - \delta Z_{ii}. \quad (B.23)
\]
where \(i = 2, 3, 4\) and the selfenergy function of Majorana particles is decomposed as
\[
\Sigma_f (\hat{q}) \equiv \Sigma^S_f (q^2) \mathbf{1} + \Sigma^V_f (q^2) \hat{q}. \quad (B.24)
\]
The neutralino wavefunction renormalization constant appearing in (B.23) is expressed in terms of the basic renormalization constants introduced in section 2 as
\[
\delta Z_{ii} = \sum_k (Re \delta Z)_k (\mathcal{O}_N)_{ik} (\mathcal{O}_N)_{ik}, \quad \text{no sum over } i, \quad (B.25)
\]
where
\[
\delta Z_k \equiv (\delta Z_{\lambda}, \delta Z_{\lambda^0}, \delta Z_{\tilde{H}_1}, \delta Z_{\tilde{H}_2}). \quad (B.26)
\]
and $\mathcal{O}$ is the orthogonal matrix which diagonalizes the neutralino mass matrix. See (A.37).

Note that for unstable particles, even if the residue condition is imposed on the propagator at the pole position, there is a non-vanishing $\delta Z_{\chi_i}^{\text{ext}}$ which is ultraviolet-finite and purely imaginary. For example, we can easily check that $\delta Z_{\chi_i}^{\text{ext}}$ which is obtained from (B.21) by changing the index 2 to 1, is purely imaginary. We can neglect such contributions if perturbation works.

References

[1] For review see, for example, P. Fayet and S. Ferrara, Phys. Rep. 32 (1977) 249; H. Haber and G. Kane, Phys. Rep. 117 (1985) 75.

[2] H.P. Nilles, Phys. Rep. 110 (1984) 1.

[3] S. L. Glashow, Nucl. Phys. 22 (1961) 579, S. Weinberg, Phys. Rev. Lett. 19 (1967) 1264, A. Salam, Proceedings of the Nobel symposium, 1968, Lerum, Sweden.

[4] J. Fujimoto et al., Comp. Phys. Comm. 111 (1998) 185.

[5] J. Fujimoto et al., Comp. Phys. Comm. 153 (2003) 106.

[6] T. Ishikawa et al., KEK Report 92-19, 1993, The GRACE manual v1.0; F. Yuasa et al., Prog. Theor. Phys. Suppl. 138 (2000) 18, [hep-ph/0007053].

[7] J. Fujimoto, T. Ishikawa, M. Jimbo, T. Kon and M. Kuroda, Proceedings of the XVIth International Workshop on High Energy Physics and Quantum Field Theory, (2004) 26, [hep-ph/0402144].

[8] J. Fujimoto, T. Ishikawa, M. Jimbo, T. Kon and M. Kuroda, Nucl. Instrum. Methods Phys. Res. A 534 (2004) 246

[9] J. Fujimoto et al., Nucl. Phys. Proc. Suppl. 157 (2006) 157.

[10] M. Kuroda, in the Research report to the Ministry of Education, Science and Culture, Japan, the Grant-in-aid for scientific research C(No. 08640391) (1999) p.127.

27
[11] M. Kuroda, KEK-CP 080 (1999), [hep-ph/9902340].

[12] K. Hikasa, SUSY manuscript version July 5, 1995, (1995), [unpublished].

[13] J. Rosiek, Phys. Rev. D41 (1990) 3464; erratum KA-TA-8-1995.

[14] P. H. Chankowski, S. Pokorski and J. Rosiek, Nucl. Phys. B423 (1994) 437.

[15] Y. Yamada, Phys. Lett. B530 (2002) 174.

[16] J. Guasch, J. Sola and W. Hollik, Phys. Lett. B437 (1998) 88.

[17] A. Yamada, Phys. Lett. B263 (1991) 233; Z. Phys. C61 (1994) 247.

[18] A. Dabelstein, Z. Phys. C67 (1995) 495.

[19] D. Pierce and A. Papadopoulos, Phys. Rev. D50 (1994) 565; Nucl. Phys. B430 (1994) 278.

[20] S. Alam, K. Hagiwara, S. Kanemura, R. Szalapski and Y. Umeda, Phys. Rev. D62 (2000) 095011.

[21] W. Hollik, E. Kraus and M. Roth, C. Rupp, K. Sibold, D. Stöckinger, [hep-ph/0204350v2].

[22] J. Guasch, W. Hollik, J. Sola, JHEP 0210 (2002) 040, [hep-ph/0207364].

[23] T. Fritzscbe and W. Hollik, Eur. Phys. J. C24 (2002) 619.

[24] G. Bélanter et al., Phys. Rep. 430 (2006) 117.

[25] J.A. Aguilar-Saaedra et al., Eur. Phys. J. C46 (2006) 43.

[26] S. Heinemeyer, W. Hollik, J. Rosiek and G. Weiglein, Eur. Phys. J. C19 (2001) 535.

[27] W. Öller, H. Eberl and W. Majerotto, Phys. Rev. D71 (2005) 115002.