# Two-loop $\mathcal{O}(\alpha_s^2)$ MSSM corrections to the pole masses of heavy quarks

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**Abstract**

We present results for two-loop MSSM corrections to the relation between pole and running masses of heavy quarks up to the order $\mathcal{O}(\alpha_s^2)$. The running masses are defined in $\overline{\mathrm{DR}}$ scheme, usually used in multiloop calculations in supersymmetric theories. An analysis of the value of these corrections in different regions of the parameter space is provided.

# 1 Introduction

Recently a lot of theoretical and experimental efforts were spent to find some new physics beyond that described by Standard Model (SM). One of the most popular ways to extend the SM is to make our world supersymmetric at some higher scale, then those currently tested by experimental facilities over the world. A way to do it minimally and at the same time to have renormalizable local quantum filed theory is to consider $N = 1$ SUSY field theory. It is obtained from the SM by replacing SM fields with groups of fields, forming representations of $N = 1$ SUSY algebra, and introducing additional Higgs doublet, required by supersymmetry. This way we obtain the Minimal Supersymmetric Standard Model (MSSM) [1, 2, 3].

A huge theoretical research in this direction was made in last years as well as parameter space of this model was heavily constrained by experimentalists. Today people consider radiative SUSY corrections to SM precision observables and try to observe SUSY particles directly in the experiment. The precision with which masses of heavy quarks are known at present, requires to take into account in addition to leading one-loop SUSY corrections also higher order corrections.

It is the aim of present article to calculate the "leading" $\mathcal{O}(\alpha_s^2)$ MSSM corrections to the pole masses of heavy quarks and to see how they affect the SUSY particle spectra, obtained from renormalization group analysis with universal boundary conditions at GUT scale. In fact we will be interested in the relation between the pole mass of heavy quark and the running heavy quark $\overline{\mathrm{DR}}$ mass. The latter is defined as a quantity computed in dimensional reduction and renormalized minimally. [1] By the "leading" corrections we mean the corrections, which are mediated by strong interactions and formally they should be dominant. However, as it is known from one-loop calculations, in the case of $b$-quark there exists large contribution from stop-chargino loop enhanced by $t$-quark Yukawa coupling, $\tan\beta$ or by supersymmetric Higgs mass parameter $\mu$. So, in this sense, our results for $b$-quark are not complete. But, they can be used to see the typical size of these corrections and already make some conclusions on whether they should be accounted for in accurate analysis or not.

To compute needed diagrams we use the large mass expansion procedure and for the moment we restrict ourselves only by terms up to $\mathcal{O}(m_b^2/M_{\rm SUSY}^2)$ order in a case of $b$-quark and up to $\mathcal{O}(m_t^2/M_{\rm SUSY}^2)$ order in the case of $t$-quark. It is a good approximation for $b$-quark. However, as our numerical analysis shows, such approximation does not work good in the case of $t$-quark and higher terms in large mass expansion should be taken into account. We suppose to calculate missing corrections from stop-chargino loops to $b$-quark pole mass as well as to supply more terms in the expansion in a relation between top pole and $\overline{\rm DR}$ masses in one of our next papers.

The paper is organized as follows. In section 2 we define quantities we want to compute and make comments on the choice of our renormalization scheme. In section 3 we describe our model — the supersymmetric QCD (a subset of MSSM, relevant to the calculation of $\alpha_s^n$ corrections) and present in detail the renormalization procedure. Section 4 contains our results and numerical analysis of the effect, produced by our correction on supersymmetric particle spectra. The latter were computed with a help of SoftSUSY program [6] under universal boundary conditions at GUT scale. And, finally, section 5 contains our conclusion.

# 2 Pole mass and choice of renormalization scheme

First, we would like to note that there are several quark mass definitions. This is mainly due to the fact, that quarks were not ever observed as free particles. Therefore the definition of their masses rely heavily on theoretical constructions. Different definitions exist refering to different renormalization schemes used in quantum field theories. We will mention only a few most popular — $\overline{\rm MS}$ mass $\overline{m}$, pole mass $M$ and $\overline{\rm DR}$ mass $m$. In the nonsupersymmetric QCD the relations between these three masses are known (to two-loops between $\overline{\rm DR}$ and pole masses [7] and to two [8, 9] and to three [10, 11] loops between $\overline{\rm MS}$ and pole masses) and in supersymmetric QCD case exists only one-loop relation between $\overline{\rm DR}$ and pole masses of heavy quarks [12, 13, 14, 15, 16]. In this work we will establish the two-loop relation between $\overline{\rm DR}$ and pole masses in supersymmetric QCD.

In $\overline{\rm DR}$ scheme particle masses and couplings depend explicitly on renormalization scale $\mu$, which is taken in applications equal to characteristic scale of a studied process. The renormalization scale dependence of these quantities could be conveniently described by renormalization group equations. They are known very well for MSSM and supersymmetric QCD. It should be noted, that $\overline{\rm DR}$ mass is a short distance quantity. That is, it is sensitive only to short distance effects. To describe processes with the characteristic scale of the order of quark mass itself a different on-shell mass definition used and here pole mass of the particle comes into play.

Pole mass of the particle is defined by the singularity of corresponding two-point function. As explicit perturbative calculations showed, the pole mass of quark is an infrared finite and gauge invariant quantity. For this reason it is considered as a physically meaningful quantity [17]. We will

[1] $\overline{\rm DR}$ scheme [4], which is similar to $\overline{\rm MS}$ scheme, is usually used for calculation in supersymmetric models, because contrary to $\overline{\rm MS}$ scheme it manifestly preserves supersymmetry, at least for not very high loop orders [5]. The details about renormalization scheme used will be given below.

restrict ourselves to perturbation theory only and will not analyze the exact nature of the two-point function singularity which may involve some nonperturbative dynamics. In the present paper we employ the definition of the pole quark mass, where it equals to the value of $p$, at which inverse quark propagator turns to zero.

The pole mass $M$ of a quark is defined as a formal solution for $\hat{p}$ (in the Minkowski metric) at which the reciprocal of the connected full propagator equals to zero:

$$\hat{p} - m - \Sigma(\hat{p}, m) = 0, \tag{1}$$

where $\Sigma(\hat{p}, m) = m\Sigma_1(p^2, m) + (\hat{p} - m)\Sigma_2(p^2, m)$ is the one-particle-irreducible two-point function, $m$ may stand for the bare or renormalized mass, $m_{\text{bare}}$ or $m_{\text{ren}}$, depending on the prescription used in evaluating $\Sigma$. The solution to eq. (1) is sought order by order in the perturbation theory. To two loops we have

$$M = m + \Sigma^{(1)}(m, m) + \Sigma^{(2)}(m, m) + \Sigma^{(1)}(m, m)\ \Sigma^{(1)'}(m, m) + \mathcal{O}(\Sigma^{(3)}), \tag{2}$$

where $\Sigma^{(L)}$ is the $L$-loop contribution to $\Sigma$, and the prime denotes the derivative with respect to the first argument. In what follows, we will be interested in the relation between pole quark mass $M$ and running $\overline{\text{DR}}$ mass $m$, computed in MSSM up to $\mathcal{O}(\alpha_s^2)$ order. Technically, to solve this problem, we need to evaluate $\approx 200$ two-loop propagator type diagrams involving many different mass scales (see Fig. 1). In general, it is quite complex problem to be solved exactly. However, in our case all mass scales can be divided in two groups denoted by $m_{\text{soft}}$ and $m_{\text{hard}}$, such that scales from $m_{\text{soft}}$ much less then each of the scales from $m_{\text{hard}}$ group:

| | $m_{\text{soft}}$ | $m_{\text{hard}}$ |
|---|---|---|
| case of $b$-quark: | $m_b$ | $m_t$, $m_{\tilde{q}}$, $m_{\tilde{g}}$ |
| case of $t$-quark: | $m_t$ | $m_{\tilde{q}}$, $m_{\tilde{g}}$ |

Here $m_{\tilde{g}}$ is gluino mass, $m_{\tilde{q}}$ stands for different squark masses. We also explicitly denoted which mass scales belong to soft or hard groups in the cases of $b$- and $t$-quarks.

Given this mass hierarchy one can employ the large mass expansion procedure to reduce evaluation of multi-scale two-loop integrals to the calculation of single scale on-shell two-loop integrals, two-loop tadpole integrals with two scales and products of one-loop on-shell integrals and one-loop tadpole diagrams. To compute two-loop single scale on-shell integrals we made use of ONSHELL2 package [18]. For the evaluation of two-loop tadpole diagrams with two scales the recurrence relations of [19] were used. Calculation of one-loop self-energies, including their derivatives with respect to momentum is an easy task and we will not make further comments on it.

Now let us make some comments on the choice of the renormalization scheme. Here we use the same renormalization prescription, as that used by authors of [7]. To be more specific, we use regularization by dimensional reduction, a modification of the conventional dimensional regularization, originally proposed in [4]. In this procedure the vector and spinor algebra in the numerator of Feynman diagrams are four-dimensional, which is the requirement imposed by supersymmetry. However, to make sense from divergent momentum integrals we need to regularize them by non-integer space-time dimension $d = 4 - 2\varepsilon$. In order for quantum corrections not to break gauge invariance we need a cancellation of squared momenta both in numerator and denominator of Feynman integrals. Thus, the momenta should form a $d$-dimensional subspace in four dimensions. As the momenta become $d$-dimensional, four-vectors naturally split into true $d$-vectors and so-called $\varepsilon$-scalars, obtained in the process of dimensional reduction of original four-dimensional lagrangian. The $\varepsilon$-scalars is nothing, but matter fields. Their appearance is the only difference from the conventional dimensional regularization.

We would like to note, that in general renormalization of $\varepsilon$-scalars and their interactions is not identical to that of vectors and so the original four-covariance may be spoiled by quantum corrections. Moreover, quantum corrections may also generate a mass for $\varepsilon$-scalars. There is an arbitrariness in choosing the renormalization scheme for this mass [20]. A consistent way is to choose the finite $\varepsilon$-scalar mass counterterm so that the pole (and renormalized) mass of the $\varepsilon$-scalars equals zero. The $\varepsilon$-scalar field renormalization is left minimal. Other renormalizations are done minimally, that is by subtracting only poles in $\varepsilon$.

# 3 Definition of the model and renormalization procedure

For our purposes here we do not need the complete MSSM but only its part — supersymmetric extension of QCD (SUSY QCD). In the superfield formulation SUSY QCD lagrangian consists of two parts: the rigid supersymmetric part and the part containing soft supersymmetry breaking terms.

$$\begin{aligned}\mathcal{L}_{\rm rigid} &= \int d^2\theta\, \frac{1}{2}\,{\rm Tr} W^\alpha W_\alpha + \int d^2\bar\theta\, \frac{1}{2}\,{\rm Tr} \bar W_{\dot\alpha} \bar W^{\dot\alpha} + \int d^2\theta\, \mathcal{W} + \int d^2\bar\theta\, \bar{\mathcal{W}} \\ &+ \int d^2\theta\, d^2\bar\theta\, \left[\bar Q e^{2g_3 T^a V^a} Q + \bar U e^{-2g_3 (T^a)^* V^a} U + \bar D e^{-2g_3 (T^a)^* V^a} D\right], \end{aligned} \tag{3}$$

where $Q = Q_L$, $U = U_R$, $D = D_R$, $T^a = \lambda^a/2$ ($a$=1,...,8) with $\lambda^a$ being the Gell-Mann matrices [2] and the gauge field strength tensors are

$$W_\alpha = -\frac{1}{8g_3}\bar D^2 e^{-2g_3 T^a V^a} D_\alpha e^{2g_3 T^a V^a}, \quad \bar W_{\dot\alpha} = -\frac{1}{8g_3} D^2 e^{2g_3 T^a V^a} \bar D_{\dot\alpha} e^{-2g_3 T^a V^a}$$

The superpotential of the ordinary SUSY QCD contains only quadratic mass terms for chiral (antichiral) superfields. Considering SUSY QCD as a part of MSSM, where all masses of particles are generated through the Higgs mechanism, the superpotential takes the following form ($\epsilon^{12} = \epsilon_{21} = 1$)

$$\mathcal{W} = \epsilon^{ij}\left[y_d D H_1^i Q^j + y_u U H_2^j Q^i + \mu H_1^j H_2^i\right], \tag{4}$$

where $y_d$ and $y_u$ are Yukawa coupling constants carrying the generation indices, which have been suppressed (as well as group indices). In this work we will need the Higgs fields only to generate masses of our particles and will not be interested in the detail structure of MSSM Higgs sector.

To perform the supersymmetry breaking, that satisfies the requirement of "softness", we introduce a gluino mass term, soft masses of scalar superpartners of quarks and soft trilinear couplings with Higgs fields, which may be written in terms of $N = 1$ superfields, provided one introduces external spurion superfields [21]

$$\begin{aligned}\mathcal{L}_{\rm soft} &= -\frac{1}{2}\int d^2\theta\, 2M_3\theta^2 {\rm Tr} W^\alpha W_\alpha - \frac{1}{2}\int d^2\bar\theta\, 2\bar M_3 \bar\theta^2 {\rm Tr} \bar W_{\dot\alpha} \bar W^{\dot\alpha} \\ &- \int d^2\theta\, d^2\bar\theta \left[M_{\tilde q}^2 \bar Q e^{2g_3 T^a V^a} Q + M_{\tilde u}^2 \bar U e^{-2g_3 (T^a)^* V^a} U + M_{\tilde d}^2 \bar D e^{-2g_3 (T^a)^* V^a} D\right] \theta^2 \bar\theta^2 \\ &- \int d^2\theta\, \epsilon^{ij}\left[y_d A_d D H_1^i Q^j + y_u A_u U H_2^j Q^i\right]\theta^2 + {\rm h.c.} \end{aligned} \tag{5}$$

[2] Note that U and D are colour anti-triplets and the anti-colour generator is $-(T^a)^*$

The chiral superfields have the following component fields representations

$$
\begin{aligned}
Q &= \tilde{q}_L - i\,\theta\sigma^\mu\bar{\theta}\partial_\mu\tilde{q}_L - \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\partial^\mu\partial_\mu\tilde{q}_L + \sqrt{2}\theta q_L - \frac{i}{\sqrt{2}}\theta\theta\bar{\theta}\bar{\sigma}^\mu\partial_\mu q_L + \theta\theta F_q\,, && (6)\\
U &= \bar{\tilde{u}}_R - i\,\theta\sigma^\mu\bar{\theta}\partial_\mu\bar{\tilde{u}}_R - \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\partial^\mu\partial_\mu\bar{\tilde{u}}_R + \sqrt{2}\theta u_R - \frac{i}{\sqrt{2}}\theta\theta\bar{\theta}\bar{\sigma}^\mu\partial_\mu u_R + \theta\theta F_u\,, && (7)\\
D &= \bar{\tilde{d}}_R - i\,\theta\sigma^\mu\bar{\theta}\partial_\mu\bar{\tilde{d}}_R - \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\partial^\mu\partial_\mu\bar{\tilde{d}}_R + \sqrt{2}\theta d_R - \frac{i}{\sqrt{2}}\theta\theta\bar{\theta}\bar{\sigma}^\mu\partial_\mu d_R + \theta\theta F_d && (8)
\end{aligned}
$$

and for the vector superfields in Wess-Zumino gauge we have

$$
V^a = \theta\sigma^\mu\bar{\theta}G^a_\mu + \theta\theta\bar{\theta}\bar{\tilde{g}}^a + \bar{\theta}\bar{\theta}\theta\tilde{g}^a + \frac{1}{2}\theta\theta\bar{\theta}\bar{\theta}D^a. \qquad (9)
$$

Then, rewriting everything in the component fields, the lagrangian of SUSY QCD takes the form ($q_R = (u_R, d_R)$)

$$
\begin{aligned}
\mathcal{L} &= -\frac{1}{4}G^a_{\mu\nu}G^{a\,\mu\nu} + i\bar{\tilde{g}}^a\bar{\sigma}^\mu D_\mu\tilde{g}^a + i\bar{q}_R\bar{\sigma}^\mu D_\mu q_R + i\bar{q}_L\bar{\sigma}^\mu D_\mu q_L\\
&+ (D^\mu\tilde{q}_R)^\dagger D_\mu\tilde{q}_R + (D^\mu\tilde{q}_L)^\dagger D_\mu\tilde{q}_L - g_3\bar{\tilde{q}}_R T^a D^a\tilde{q}_R + g_3\bar{\tilde{q}}_L T^a D^a\tilde{q}_L\\
&+ \sqrt{2}g_3 q_R T^a\tilde{g}^a\tilde{q}_R - \sqrt{2}g_3\bar{q}_L T^a\bar{\tilde{g}}^a\tilde{q}_L + \sqrt{2}g_3\bar{\tilde{q}}_R T^a\bar{\tilde{g}}^a\bar{q}_R - \sqrt{2}g_3\bar{\tilde{q}}_L T^a\tilde{g}^a q_L\\
&+ \frac{1}{2}D^a D^a + \frac{1}{2}D^i D^i + \frac{1}{2}D' D' + g_2\bar{\tilde{q}}_L\frac{\sigma^i}{2}D^i\tilde{q}_L + g_2\bar{h}_j\frac{\sigma^i}{2}D^i h_j + g_1\bar{h}^0_j\frac{Y_{h_j}}{2}D' h^0_j\\
&+ g_1\bar{\tilde{q}}_L\frac{Y_{\tilde{q}_L}}{2}D'\tilde{q}_L - g_1\bar{\tilde{q}}_R\frac{Y_{\tilde{q}_R}}{2}D'\tilde{q}_R + \bar{F}_{h^0_1}F_{h^0_1} + \bar{F}_{h^0_2}F_{h^0_2} + \bar{F}_{q_L}F_{q_L} + \bar{F}_{q_R}F_{q_R}\\
&- \Big[\mu h^0_1 F_{h^0_2} + \mu F_{h^0_1}h^0_2 - y_d\bar{\tilde{d}}_R F_{h^0_1}\tilde{d}_L - y_u\bar{\tilde{u}}_R F_{h^0_2}\tilde{u}_L\\
&- y_d(F_{d_R}h^0_1\tilde{d}_L + \bar{\tilde{d}}_R h^0_1 F_{d_L}) - y_u(F_{u_R}h^0_2\tilde{u}_L + \bar{\tilde{u}}_R h^0_2 F_{u_L})\\
&+ y_d d_R h^0_1 d_L + y_u u_R h^0_2 u_L + y_d A_d\tilde{d}_R h^0_1\tilde{d}_L + y_u A_u\tilde{u}_R h^0_2\tilde{u}_L + \text{h.c.}\Big]\\
&- \frac{M_3}{2}\tilde{g}^a\tilde{g}^a - \frac{M_3}{2}\bar{\tilde{g}}^a\bar{\tilde{g}}^a - M^2_{\tilde{q}}\bar{\tilde{q}}_L\tilde{q}_L - M^2_{\tilde{u}}\bar{\tilde{u}}_R\tilde{u}_R - M^2_{\tilde{d}}\bar{\tilde{d}}_R\tilde{d}_R,
\end{aligned}
$$

where gauge covariant derivative and field strength are defined as

$$
\begin{aligned}
D_\mu &= \partial_\mu + ig_3 T^a G^a_\mu\,,\\
G^a_{\mu\nu} &= \partial_\mu G^a_\nu - \partial_\nu G^a_\mu - g_3 f^{abc}G^b_\mu G^c_\nu\,.
\end{aligned}
$$

Here we also added auxiliary $SU(2)\times U(1)$ fields $D^i$ and $D'$, which are necessary to describe correctly the mass generation mechanism. $\sigma^i$ denote Pauli matrices and $Y_q$ is a hypercharge of a particle $q$. Eliminating auxiliary fields $F_q$, $\bar{F}_q$, $D^a$, $D^i$ and $D'$ and introducing four-component spinor notation similar to [1, 22] one gets

$$
\begin{aligned}
\mathcal{L} &= -\frac{1}{4}G^a_{\mu\nu}G^{a\,\mu\nu} + \frac{i}{2}\bar{\tilde{g}}^a\gamma^\mu D_\mu\tilde{g}^a + i\bar{q}\gamma^\mu D_\mu q + (D^\mu\tilde{q}_R)^\dagger D_\mu\tilde{q}_R + (D^\mu\tilde{q}_L)^\dagger D_\mu\tilde{q}_L\\
&+ \sqrt{2}g_3\bar{q}P_L T^a\tilde{g}^a\tilde{q}_R - \sqrt{2}g_3\bar{q}P_R T^a\tilde{g}^a\tilde{q}_L + \sqrt{2}g_3\bar{\tilde{q}}_R T^a\bar{\tilde{g}}^a P_R q - \sqrt{2}g_3\bar{\tilde{q}}_L T^a\bar{\tilde{g}}^a P_L q\\
&- \frac{g_3^2}{2}(\bar{\tilde{q}}_L T^a\tilde{q}_L - \bar{\tilde{q}}_R T^a\tilde{q}_R)(\bar{\tilde{q}}_L T^a\tilde{q}_L - \bar{\tilde{q}}_R T^a\tilde{q}_R) - \frac{M_3}{2}\bar{\tilde{g}}^a\tilde{g}^a + \mathcal{L}_M\,, \qquad (10)\\
\mathcal{L}_M &= -M^2_{\tilde{q}}\bar{\tilde{q}}_L\tilde{q}_L - M^2_{\tilde{u}}\bar{\tilde{u}}_R\tilde{u}_R - M^2_{\tilde{d}}\bar{\tilde{d}}_R\tilde{d}_R - y_d h^0_1\bar{d}d - y_u h^0_2\bar{u}u
\end{aligned}
$$

$$
\begin{aligned}
&- \left(\frac{g_2^2}{2} I^3_{\tilde{q}_L} \bar{\tilde{q}}_L \tilde{q}_L - \frac{g_1^2}{4} Y_{\tilde{q}_L} \bar{\tilde{q}}_L \tilde{q}_L + \frac{g_1^2}{4} Y_{\tilde{q}_R} \bar{\tilde{q}}_R \tilde{q}_R \right) \left(\bar{h}_1^0 h_1^0 - \bar{h}_2^0 h_2^0\right) \\
&- y_d^2 \left(h_1^0 \tilde{d}_L\right)^\dagger h_1^0 \tilde{d}_L - y_d^2 \bar{h}_1^0 h_1^0 \bar{\tilde{d}}_R \tilde{d}_R - y_u^2 \left(h_2^0 \tilde{u}_L\right)^\dagger h_2^0 \tilde{u}_L - y_u^2 \bar{h}_2^0 h_2^0 \bar{\tilde{u}}_R \tilde{u}_R \\
&- \left[ y_d h_1^0 A_d \bar{\tilde{d}}_R \tilde{d}_L + y_u h_2^0 A_u \bar{\tilde{u}}_R \tilde{u}_L - y_d \mu h_2^0 \bar{\tilde{d}}_R \tilde{d}_L - y_u \mu h_1^0 \bar{\tilde{u}}_R \tilde{u}_L + \text{h.c} \right] \qquad (11)
\end{aligned}
$$

where $P_L = (1-\gamma_5)/2$ and $P_R = (1+\gamma_5)/2$ are the chiral and antichiral projectors.

After Higgsing and taking into account that

$$
\begin{aligned}
& m_d = y_d v_1, \qquad m_u = y_u v_2, \qquad \tan\beta = \frac{v_2}{v_1}, \qquad \sin^2\theta_W = \frac{g_1^2}{g_1^2+g_2^2}, \\
& m_Z^2 = \frac{1}{2}(g_1^2+g_2^2)(v_1^2+v_2^2) = \frac{1}{2}\frac{g_1^2}{\sin^2\theta_W}(v_1^2+v_2^2), \\
& \cos 2\beta = \frac{v_1^2-v_2^2}{v_1^2+v_2^2}, \qquad e_q = \frac{Y_q}{2} + I_q^3,
\end{aligned}
$$

it is convenient to rewrite $\mathcal{L}_M$ from eq. (11) in the following form [23]

$$
\mathcal{L}_M = -m_q \bar{q} q - m^2_{\tilde{q}_L} \bar{\tilde{q}}_L \tilde{q}_L - m^2_{\tilde{u}_R} \bar{\tilde{u}}_R \tilde{u}_R - m^2_{\tilde{d}_R} \bar{\tilde{d}}_R \tilde{d}_R - a_q m_q \bar{\tilde{q}}_L \tilde{q}_R - a_q m_q \bar{\tilde{q}}_R \tilde{q}_L \qquad (12)
$$

with

$$
\begin{aligned}
m^2_{\tilde{q}_L} &= M^2_{\tilde{q}} + m_Z^2 \cos 2\beta \, (I^3_{\tilde{q}_L} - e_q \sin^2\theta_W) + m_q^2 \,, \\
m^2_{\tilde{u}_R} &= M^2_{\tilde{u}} + e_u \, m_Z^2 \cos 2\beta \, \sin^2\theta_W + m_u^2 \,, \qquad a_u = A_u - \mu \cot\beta \\
m^2_{\tilde{d}_R} &= M^2_{\tilde{d}} + e_d \, m_Z^2 \cos 2\beta \, \sin^2\theta_W + m_d^2 \,, \qquad a_d = A_d - \mu \tan\beta
\end{aligned}
$$

where $e_q$ and $I_q^3$ are the electric charge and the third component of the weak isospin of a particle $q$, $m_q$ is the mass of the partner quark and $\mu$ is a Higgs mass parameter. $M_{\tilde{q}}$, $M_{\tilde{u}}$, and $M_{\tilde{d}}$ are soft SUSY breaking masses, $A_u$ and $A_d$ are trilinear couplings as in eq. (5). Family indices have been suppressed.

Now we must diagonalize the mass matrix $\mathcal{M}^2_{\tilde{q}}$ of squarks to determine the physical mass eigenstates

$$
\mathcal{M}^2_{\tilde{q}} = \begin{pmatrix} m^2_{\tilde{q}_L} & a_q m_q \\ a_q m_q & m^2_{\tilde{q}_R} \end{pmatrix} = (\mathcal{R}^{\tilde{q}})^\dagger \begin{pmatrix} m^2_{\tilde{q}_1} & 0 \\ 0 & m^2_{\tilde{q}_2} \end{pmatrix} \mathcal{R}^{\tilde{q}} \,. \qquad (13)
$$

According to eq. (13) $\mathcal{M}^2_{\tilde{q}}$ is diagonalized by an unitary matrix $\mathcal{R}^{\tilde{q}}$. Assuming that CP violating phases occur only in the CKM matrix, we choose $\mathcal{R}^{\tilde{q}}$ to be real. The weak eigenstates $\tilde{q}_L$ and $\tilde{q}_R$ are related thus to their mass eigenstates $\tilde{q}_1$ and $\tilde{q}_2$ by

$$
\begin{pmatrix} \tilde{q}_1 \\ \tilde{q}_2 \end{pmatrix} = \mathcal{R}^{\tilde{q}} \begin{pmatrix} \tilde{q}_L \\ \tilde{q}_R \end{pmatrix}, \qquad \mathcal{R}^{\tilde{q}} = \begin{pmatrix} \cos\theta_{\tilde{q}} & \sin\theta_{\tilde{q}} \\ -\sin\theta_{\tilde{q}} & \cos\theta_{\tilde{q}} \end{pmatrix}, \qquad (14)
$$

with $\theta_{\tilde{q}}$ being the squark mixing angle. The mass eigenvalues are given by

$$
m^2_{\tilde{q}_{1,2}} = \frac{1}{2}\left( m^2_{\tilde{q}_L} + m^2_{\tilde{q}_R} \mp \sqrt{(m^2_{\tilde{q}_L} - m^2_{\tilde{q}_R})^2 + 4\, a_q^2 m_q^2} \right). \qquad (15)
$$

By convention, we choose $\tilde{q}_1$ to be the lightest mass eigenstate. Notice, that $m_{\tilde{q}_1} \le m_{\tilde{q}_{L,R}} \le m_{\tilde{q}_2}$.

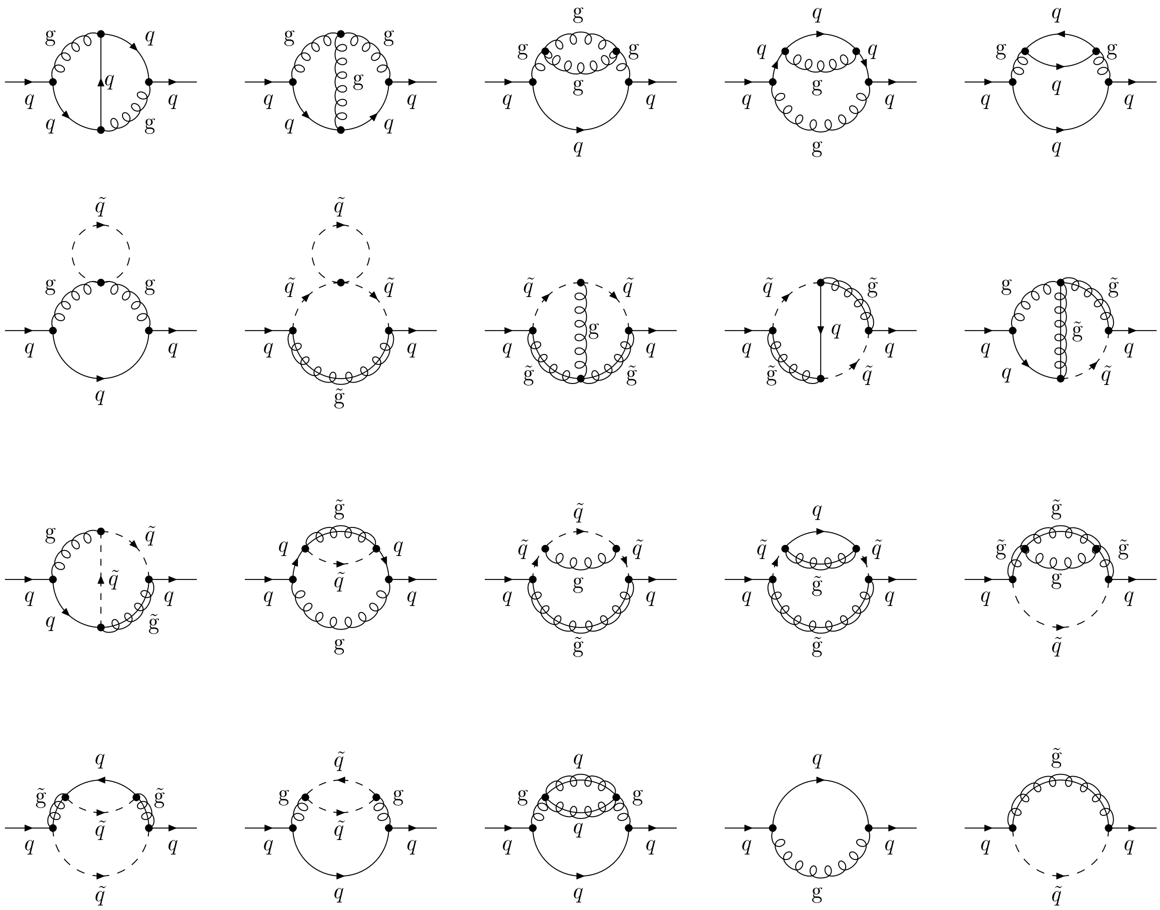


Figure 1: Feynman diagrams for $\mathcal{O}(\alpha_s^2)$ contribution to quark self-energies (last two diagrams at the one loop level of order $\mathcal{O}(\alpha_s)$). The diagrams in the first line are the pure QCD contribution.

For the mixing angle $\theta_{\tilde{q}}$ we require $0 \le \theta_{\tilde{q}} < \pi$. Then one can write

$$\cos\theta_{\tilde{q}} = \frac{-a_q m_q}{\sqrt{(m_{\tilde{q}_L}^2 - m_{\tilde{q}_1}^2)^2 + a_q^2 m_q^2}}, \qquad \sin\theta_{\tilde{q}} = \frac{m_{\tilde{q}_L}^2 - m_{\tilde{q}_1}^2}{\sqrt{(m_{\tilde{q}_L}^2 - m_{\tilde{q}_1}^2)^2 + a_q^2 m_q^2}}. \tag{16}$$

Moreover, $|\cos\theta_{\tilde{q}}| > \frac{1}{\sqrt{2}}$ if $m_{\tilde{q}_L} < m_{\tilde{q}_R}$ and $|\cos\theta_{\tilde{q}}| < \frac{1}{\sqrt{2}}$ if $m_{\tilde{q}_R} < m_{\tilde{q}_L}$.

Making transition in eq. (10) from $\tilde{q}_R$ and $\tilde{q}_L$ to $\tilde{q}_1$ and $\tilde{q}_2$ with the help of the matrix from eq. (14) we get for quark-gluino-squark trilinear and four-squarks vertices complicated expressions due to squark mixing, which we do not present here. To produce input amplitudes of the diagrams needed in our calculations we made use of Mathematica package FeynArts [24], which have SUSY QCD as part of MSSM model [25]. The Feynman diagrams contributing to quarks self-energies in one and two loop order are shown in Fig. 1.

Note again, that doing these calculations we work in the minimal subtraction scheme $\overline{\mathrm{MS}}$ adopted to supersymmetry — dimensional reduction with minimal subtraction $\overline{\mathrm{DR}}$ . To obtain a relation between pole and $\overline{\mathrm{DR}}$ quark masses we need to know $\overline{\mathrm{DR}}$ renormalization constants for quark, gluino, squark, $\varepsilon$-scalar masses, gauge and Yukawa charges[3] as well as renormalization of squark mixing angle. We will describe evaluation of each of these quantities in next subsections.

[3]By Yukawa charge we mean the Yukawa coupling of $\varepsilon$-scalar to fermions

### 3.1 Renormalization of QCD sector

First we have to renormalize the charge and the masses of quarks as well as fields. The relations between bare and renormalized parameters are given by

$$\begin{aligned}
\alpha_{s,0} &= \mu^{\varepsilon/2} Z_{\alpha_s} \alpha_s \,, &(17)\\
m_{q,0} &= Z_{m_q} m_q \,, &(18)\\
G_0 &= Z_G^{1/2} G \,, &(19)\\
q_0 &= Z_q^{1/2} q \,, &(20)
\end{aligned}$$

where $\mu$ is a renormalization scale.

The charge renormalization constant can be obtained from the renormalization constants of the fields and vertex with the use of the following relation:

$$Z_{\alpha_s} = Z_{qGq}^2 \, Z_q^{-2} \, Z_G^{-1} \,, \tag{21}$$

where $Z_q$ and $Z_G$ correspond to renormalizations of quark and gluon fields and $Z_{qGq}$ renormalizes the quark-gluon interaction vertex.

One can also renormalize the gauge parameter[4] with the help of introduced above constant $Z_G$

$$\xi_0 = Z_G \xi \,. \tag{22}$$

However in the gauge invariant quantities (like e.g. the pole mass) the gauge parameter drops out already from the bare expression and renormalization (22) is not needed.

Our results for renormalization constants in $\overline{\mathrm{DR}}$ scheme read

$$\begin{aligned}
Z_{m_q} &= 1 - C_F \frac{\alpha_s}{2\pi} \frac{1}{\varepsilon} - C_F \left(\frac{\alpha_s}{4\pi}\right)^2 \left( (6 - 3C_A - 2C_F) \frac{1}{\varepsilon^2} + (-6 + 3C_A - 2C_F) \frac{1}{\varepsilon} \right), &(23)\\
Z_q &= 1 - \frac{\alpha_s}{4\pi} C_F (\xi + 1) \frac{1}{\varepsilon}, &(24)\\
Z_G &= 1 - \frac{\alpha_s}{4\pi} \left( 6 - \frac{3 - \xi}{2} C_A \right) \frac{1}{\varepsilon}, &(25)\\
Z_{qGq} &= 1 - \frac{\alpha_s}{4\pi} \left( (1 + \xi) C_F + \frac{3 + \xi}{4} C_A \right) \frac{1}{\varepsilon}, &(26)\\
Z_{\alpha_s} &= 1 - \frac{\alpha_s}{4\pi} (3C_A - 6) \frac{1}{\varepsilon}. &(27)
\end{aligned}$$

For $SU(N_c)$ we have $C_F = (N_c^2 - 1)/(2N_c)$ and $C_A = N_c$. Here we explicitly put the number of quark flavors to six.

There are several possibilities to check the correctness of our results. The coupling renormalization constant (27) can be compared with that, obtained from the one-loop MSSM beta function for $SU(3)$ group. We have also checked, that in the case of ordinary QCD we recover renormalization constants given in [7]. The 2-loop $\overline{\mathrm{DR}}$ quark mass renormalization constant (23) has been obtained from eq. (2) and contains poles in $\varepsilon$ remaining after all other renormalizations were done. Here we mean, that one should perform one-loop renormalizations of strong coupling constant and particle masses. A subtle point is the renormalization of Yukawa interaction of $\varepsilon$-scalars and quarks. In

[4] Gauge parameter $\xi$ is defined in such a way that $\xi = 0$ and $\xi = 1$ correspond to Landau and Feynman gauge respectively.

SUSY QCD discussed here, one do not need to introduce separate coupling constant $Y$ for this interaction, as supersymmetry insures that it renormalizes in the same way as strong coupling constant[5]. However, to make our renormalization procedure as general as possible ( and also applicable in a case of nonsupersymmetric QCD for checks) we introduce additional independent constant $Y$ and renormalize it separately. So, here we follow renormalization procedure described in [7]. Renormalizing a physical quantity, like pole mass, we ignore field renormalization constants, which is valid procedure, provided we are evaluating the contributions of counterterms for the whole sum of diagrams. $1/\varepsilon^2$ poles in $Z_{m_q}$ can be checked with the use of renormalization group technique. Indeed, let us write

$$\begin{aligned} m_0 &= m\left(1+\frac{\alpha_s}{4\pi}\frac{Z^{(1,1)}}{\varepsilon}+\left(\frac{\alpha_s}{4\pi}\right)^2\frac{Z^{(2,1)}}{\varepsilon}+\left(\frac{\alpha_s}{4\pi}\right)^2\frac{Z^{(2,2)}}{\varepsilon^2}\right) \\ &= m\left(1+\frac{z^{(1)}}{\varepsilon}+\frac{z^{(2)}}{\varepsilon^2}+\ldots\right), \end{aligned} \tag{28}$$

where $z^{(n)}=\sum_{l=1}^{\infty}(\alpha_s/4\pi)^l Z^{(l,n)}$. Now using renormalization group equation[6]

$$0=\mu^2\frac{d}{d\mu^2}m_0$$

we obtain

$$\begin{aligned} \gamma &= \frac{1}{2}g\frac{\partial}{\partial g}z^{(1)}, \\ \left(\gamma+\beta_g\frac{\partial}{\partial g}+\sum_i\gamma_i m_i\frac{\partial}{\partial m_i}\right)z^{(n)} &= \frac{1}{2}g\frac{\partial}{\partial g}z^{(n+1)}, \end{aligned} \tag{29}$$

($g$ is related to $\alpha_s$ via $g^2/4\pi=\alpha_s$). The formula presented are general and also applicable in a theory with spontaneously broken symmetry [36]. In the case of SUSY QCD from the previous formula it follows that

$$2Z^{(2,2)}=\left(Z^{(1,1)}\right)^2+2\beta_g^{(1)}Z^{(1,1)}. \tag{30}$$

Substituting into this relation the expressions for $Z^{(1,1)}=-2C_F$ and $\beta_g^{(1)}=-\frac{1}{2}(3C_A-6)$ we get the same result for $1/\varepsilon^2$ poles as in eq. (23).

## 3.2 Gluino mass renormalization

The renormalization of gluino mass is performed in a way similar to quark mass renormalization. Relevant diagrams with gluino-gluon and quark-squark loops are shown in Fig.2. Explicit calculations give the following results for gluino mass $m_{\tilde{g},0}=Z_{m_{\tilde{g}}}m_{\tilde{g}}$, gluino wave function renormalization

[5] See section about $\varepsilon$-scalar sector renormalization for details

[6] In fact $m_0$ is not renormalization group invariant in our renormalization scheme. However, to check leading poles in $\varepsilon$ we can treat it as RG invariant, as different prescriptions used to renormalize $\varepsilon$-scalar sector give different results only at subleading order in $\varepsilon$. This point will be discussed in detail in section concerning $\varepsilon$-scalar sector renormalization.

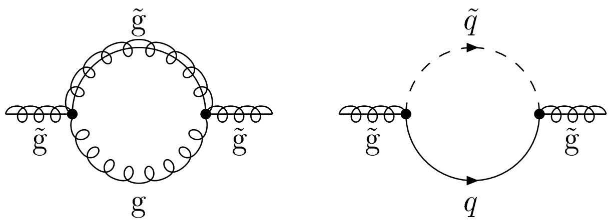

Figure 2: One-loop diagrams contributing to gluino mass counterterm.

constants $Z_{\tilde{g}}$ as well as for gluino-gluon-gluino vertex renormalization constant $Z_{\tilde{g}G\tilde{q}}$

$$
\begin{aligned}
Z_{m_{\tilde{g}}} &= 1 + \frac{\alpha_s}{4\pi}\left(6 - 3C_A\right)\frac{1}{\varepsilon}, && (31)\\
Z_{\tilde{g}} &= 1 - \frac{\alpha_s}{4\pi}\left(\xi C_A + 6\right)\frac{1}{\varepsilon}, && (32)\\
Z_{\tilde{g}G\tilde{g}} &= 1 - \frac{\alpha_s}{4\pi}\left(6 + \frac{3+5\xi}{4}C_A\right)\frac{1}{\varepsilon},. && (33)
\end{aligned}
$$

From eqs. (25), (32) and (33) one can extract the gauge coupling renormalization constant

$$
Z_{\alpha_s} = Z^2_{\tilde{g}G\tilde{g}} Z^{-2}_{\tilde{g}} Z^{-1}_{G}, \qquad (34)
$$

which, of course, coincide with eq. (27). This serves as an additional check. The result for gluino mass renormalization constant can be verified also with help of known gluino mass beta-function.

## 3.3 Squark sector renormalization

To discuss the renormalization of squark sector let us start with the lagrangian written in terms of physical states. The bare lagrangian in terms of bare parameters and bare fields is given by [7]

$$
\begin{aligned}
\mathcal{L}^{\text{bare}}_{\tilde{q}} &= \partial_\mu \bar{\tilde{q}}_i \partial_\mu \tilde{q}_i - m^2_{\tilde{q}_i} \bar{\tilde{q}}_i \tilde{q}_i + \mathcal{L}^{\text{bare}}_{q\tilde{g}\tilde{q}}, && (35)\\
\mathcal{L}^{\text{bare}}_{q\tilde{g}\tilde{q}} &= -\sqrt{2} g_3 \bar{q}\left(P_R\cos\theta_{\tilde{q}} - P_L \sin\theta_{\tilde{q}}\right) T^a \tilde{g}^a \tilde{q}_1 \\
&+ \sqrt{2} g_3 \bar{q}\left(P_R\sin\theta_{\tilde{q}} + P_L \cos\theta_{\tilde{q}}\right) T^a \tilde{g}^a \tilde{q}_2 \\
&- \sqrt{2} g_3 \bar{\tilde{q}}_1 T^a \bar{\tilde{g}}^a \left(P_L\cos\theta_{\tilde{q}} - P_R \sin\theta_{\tilde{q}}\right) q \\
&+ \sqrt{2} g_3 \bar{\tilde{q}}_2 T^a \bar{\tilde{g}}^a \left(P_L\sin\theta_{\tilde{q}} + P_R \cos\theta_{\tilde{q}}\right) q && (36)
\end{aligned}
$$

We can rewrite this lagrangian in terms of renormalized physical quantities and counterterms

$$
\begin{aligned}
\mathcal{L}^{\text{bare}}_{\tilde{q}} &= \mathcal{L}_{\tilde{q}} + \delta\mathcal{L}_{\tilde{q}} && (37)\\
\delta\mathcal{L}_{\tilde{q}} &= \delta Z_{\tilde{q}}(\partial_\mu \bar{\tilde{q}}_i \partial_\mu \tilde{q}_i - m^2_{\tilde{q}_i}\bar{\tilde{q}}_i\tilde{q}_i) - \delta m^2_{\tilde{q}_i}\bar{\tilde{q}}_i\tilde{q}_i - \delta\tilde{m}^2_{12}(\bar{\tilde{q}}_2\tilde{q}_1 + \bar{\tilde{q}}_1\tilde{q}_2) + \delta Z_{q\tilde{g}\tilde{q}}\mathcal{L}_{q\tilde{g}\tilde{q}} && (38)
\end{aligned}
$$

where $\mathcal{L}_{\tilde{q}}$ and $\mathcal{L}_{q\tilde{g}\tilde{q}}$ have the same form as (35) and (36). It should be noticed, that this set of counterterms is sufficient to cancel all rising divergencies. The $\delta\tilde{m}^2_{12}$ counterterm is used to render non-diagonal squark self-energies $\tilde{\Sigma}_{\tilde{q}_i\tilde{q}_j}$ $(i \neq j)$ finite. In practical calculations, it is convenient to

[7] We don't include bare lagrangian for four-squark interaction because counterterms associated with the four-squark vertices will give contribution of order $\mathcal{O}(\alpha_s^3)$.

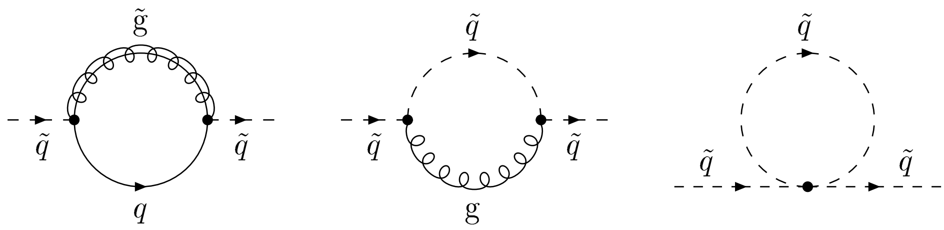


Figure 3: One-loop diagrams contributing to the squark mass and squark mixing angle counterterm.

trade $\delta\tilde{m}^2_{12}$ for the counterterm for squark mixing angle. Indeed, we can diagonalize the lagrangian (37) with the following field redefinitions

$$\begin{pmatrix} \tilde{q}'_1 \\ \tilde{q}'_2 \end{pmatrix} = \begin{pmatrix} 1 & \delta\theta_{\tilde{q}} \\ -\delta\theta_{\tilde{q}} & 1 \end{pmatrix} \begin{pmatrix} \tilde{q}_1 \\ \tilde{q}_2 \end{pmatrix} \tag{39}$$

where we have taken into account that $\delta\theta_{\tilde{q}} \sim \mathcal{O}(\alpha_s)$, $\cos(\delta\theta_{\tilde{q}}) = 1+\mathcal{O}(\alpha_s^2)$ and $\sin(\delta\theta_{\tilde{q}}) = \delta\theta_{\tilde{q}}+\mathcal{O}(\alpha_s^3)$ Demanding that the new mass matrix of squarks

$$\begin{pmatrix} 1 & \delta\theta_{\tilde{q}} \\ -\delta\theta_{\tilde{q}} & 1 \end{pmatrix} \begin{pmatrix} m^2_{\tilde{q}_1} + \delta Z_{\tilde{q}} m^2_{\tilde{q}_1} + \delta m^2_{\tilde{q}_1} & \delta\tilde{m}^2_{12} \\ \delta\tilde{m}^2_{12} & m^2_{\tilde{q}_2} + \delta Z_{\tilde{q}} m^2_{\tilde{q}_2} + \delta m^2_{\tilde{q}_2} \end{pmatrix} \begin{pmatrix} 1 & -\delta\theta_{\tilde{q}} \\ \delta\theta_{\tilde{q}} & 1 \end{pmatrix}. \tag{40}$$

must be diagonal and expanding eq. (40) in gauge coupling constant $\alpha_s$ one gets

$$\delta\theta_{\tilde{q}} = \frac{\delta\tilde{m}^2_{12}}{m^2_{\tilde{q}_1} - m^2_{\tilde{q}_2}} = \frac{(\tilde{\Sigma}^{(1)}_{\tilde{q}_1\tilde{q}_2})_{\rm div}}{m^2_{\tilde{q}_1} - m^2_{\tilde{q}_2}}, \tag{41}$$

where subscript "div" stands for the $1/\varepsilon$ part of expression. Rewriting eq. (37) in terms of new fields $\tilde{q}'_i$ we come to the following renormalization prescription for the squark sector

$$\begin{aligned} m^2_{\tilde{q},0} &= m^2_{\tilde{q}_i} + \delta m^2_{\tilde{q}_i} = Z_{m^2_{\tilde{q}_i}} m^2_{\tilde{q}_i}, && (42) \\ \tilde{q}_{i,0} &= (1 + \frac{1}{2} Z_{\tilde{q}})\tilde{q}_i = Z^{1/2}_{\tilde{q}} \tilde{q}_i && (43) \\ \theta_{\tilde{q},0} &= \theta_{\tilde{q}} + \delta\theta_{\tilde{q}} = Z_{\theta_{\tilde{q}}} \theta_{\tilde{q}} && (44) \end{aligned}$$

Explicit calculation of squark self-energies gives (see diagrams in Fig. 3) the following expression for the renormalization of squark mixing angle

$$\delta\theta_{\tilde{q}} = C_F \frac{\alpha_s}{4\pi} \left( -\frac{\sin(4\theta_{\tilde{q}})}{2} + \frac{4 m_q m_{\tilde{g}} \cos(2\theta_{\tilde{q}})}{m^2_{\tilde{q}_1} - m^2_{\tilde{q}_2}} \right) \frac{1}{\varepsilon}. \tag{45}$$

Note, that renormalization (45) of the mixing angle is minimal. It is the correct choice, if one wants to have final result expressed in terms of running $\overline{\rm DR}$ sparticle masses.

Summarizing the result of calculation the renormalization constants for the squark sector read

$$\begin{aligned} Z_{\tilde{q}} &= 1 - C_F \frac{\alpha_s}{4\pi} (\xi - 1) \frac{1}{\varepsilon}, && (46) \\ Z_{m^2_{\tilde{q}_1}} &= 1 + \frac{C_F}{m^2_{\tilde{q}_1}} \frac{\alpha_s}{4\pi} \left( \left( m^2_{\tilde{q}_2} - m^2_{\tilde{q}_1} \right) \sin^2 2\theta_{\tilde{q}} + 4 \left( m_q m_{\tilde{g}} \sin 2\theta_{\tilde{q}} - m^2_q - m^2_{\tilde{g}} \right) \right) \frac{1}{\varepsilon}, && (47) \end{aligned}$$

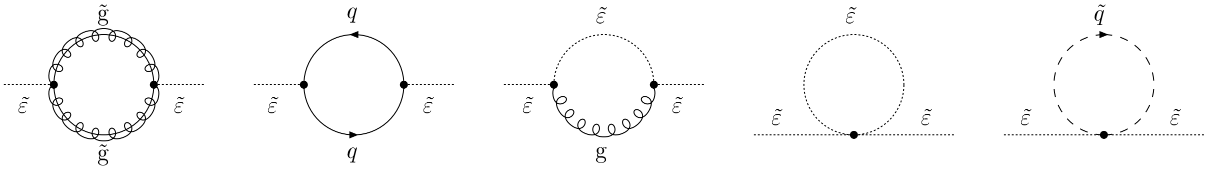


Figure 4: One-loop diagrams contributing to the $\varepsilon$-scalar wave function and mass counterterms

$$Z_{m^2_{\tilde{q}_2}} = 1 + \frac{C_F}{m^2_{\tilde{q}_2}}\frac{\alpha_s}{4\pi}\left(\left(m^2_{\tilde{q}_1} - m^2_{\tilde{q}_2}\right)\sin^2 2\theta_{\tilde{q}} - 4\left(m_q m_{\tilde{g}}\sin 2\theta_{\tilde{q}} + m_q^2 + m_{\tilde{g}}^2\right)\right)\frac{1}{\varepsilon}, \tag{48}$$

$$Z_{q\tilde{g}\tilde{q}} = 1 - \frac{\alpha_s}{4\pi}\left(\xi C_F + \frac{3+\xi}{2}C_A\right)\frac{1}{\varepsilon}, \tag{49}$$

$$Z_{\theta_{\tilde{q}}} = 1 + \frac{\delta\theta_{\tilde{q}}}{\theta_{\tilde{q}}}. \tag{50}$$

Using eqs. (24), (32), (46) and (49) one can also get gauge coupling renormalization constant $Z_{\alpha_s}$ from quark-gluino-squark vertex

$$Z_{\alpha_s} = Z^2_{q\tilde{g}\tilde{q}} Z_q^{-1} Z_{\tilde{g}}^{-1} Z_{\tilde{q}}^{-1} \tag{51}$$

which, due to supersymmetry, of cause coincides with the results of eqs. (21) and (34).

### 3.4 $\varepsilon$-scalar sector renormalization

The renormalization of $\varepsilon$-scalar sector in SUSY QCD goes along the same lines as for the ordinary QCD in dimensional reduction. Here we will remind the tricks used to obtain corresponding renormalization constants and write down their values in SUSY QCD. The part of SUSY QCD lagrangian containing $\varepsilon$-scalar fields could be obtained through the dimensional reduction of lagrangian from eq.(10). As a result we have

$$\begin{aligned}\mathcal{L}_\varepsilon = {} & \frac{1}{2}\left(D^\mu G_\sigma\right)^\dagger D_\mu G_\sigma - g_3\bar{q}\gamma_\sigma T^a G^a_\sigma q - \frac{1}{4}g_3^2 f^{abc}f^{ade}G^b_\sigma G^c_{\sigma'}G^d_\sigma G^e_{\sigma'} \\ & -\frac{1}{2}g_3\bar{\tilde{g}}\gamma_\sigma T^a G^a_\sigma \tilde{g} + g_3^2\bar{\tilde{q}}T^a T^b G^a_\sigma G^b_\sigma \tilde{q}.\end{aligned} \tag{52}$$

Here $G_{\sigma(\sigma')}$ denote $\varepsilon$-scalar fields and indices $\sigma$, $\sigma'$ belong to $2\varepsilon$ subspace.

The standard way is to use Feynman rules, derived from this lagrangian to account for contribution of $\varepsilon$-scalars. However, it turns out that for some problems one can use far more simple arguments to get a desired solution. For example, we are interested both in full and separate contributions of vectors and $\varepsilon$-scalars to the quantity, like Green function or amplitude, without external vector indices. So, vectors and $\varepsilon$-scalars fields are located on internal lines only. Then to solve this problem one should

(a) perform $d$-dimensional vector and spinor algebra

(b) for full contribution - replace $d$ in the numerator of scalarized expression with 4;
for $\varepsilon$-scalar contribution - replace $d^n$ with $4^n - (4-2\varepsilon)^n$.

When we are dealing with a Green function containing external vector indices and want to extract a contribution of either external vector or $\varepsilon$-scalar fields, a corresponding projector on $4-2\varepsilon$ or

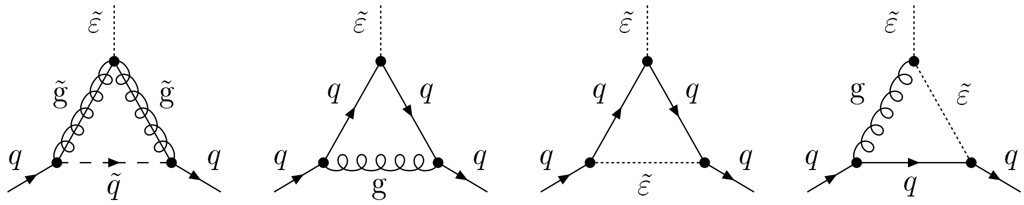


Figure 5: One-loop diagrams for $\varepsilon$-scalar Yukawa coupling renormalization constant.

$2\varepsilon$ subspaces should be constructed. Internal $\varepsilon$-scalars and vectors are treated along the same way as described above. At this point one should keep in mind, that particle momenta are always orthogonal to $2\varepsilon$ subspace. For example, in the case of $\varepsilon$-scalar propagator (see Fig.4) we use as a projector $\frac{1}{2\varepsilon}g^{\mu\nu}$ ($\varepsilon$-scalar propagator is always diagonal - see notice above) and after contraction of external indices replace $d^n$ with $4^n - (4-2\varepsilon)^n$. At the same moment, evaluating corrections to the interaction between $\varepsilon$-scalar and quarks (see Fig.5), required for the determination of Yukawa charge renormalization, one should follow the same steps as in a case of $\varepsilon$-scalar propagator with the only difference that now projector is given by

$$\mathrm{Tr}\left(\frac{1}{8C_F C_A \varepsilon} T^a \gamma^\mu \Gamma_\mu\right), \tag{53}$$

where $\Gamma^\mu$ is an interaction vertex of gluon field with quarks, computed in SUSY QCD without separating $\epsilon$-scalar contributions but using 4-dimensional vector and spinor algebras.
At the one-loop order the $\varepsilon$-scalar field renormalization constant is given by

$$Z_s = 1 - \frac{\alpha_s}{4\pi}\Big(6 - C_A(2-\xi)\Big)\frac{1}{\varepsilon}. \tag{54}$$

For the one-loop Yukawa charge renormalization constant, describing the interaction of $\epsilon$-scalar with fermions in SUSY QCD we have

$$Z_Y = Z_{qsq}^2 Z_q^{-2} Z_s^{-1} = 1 - \frac{\alpha_s}{4\pi}\left(3C_A - 6\right)\frac{1}{\varepsilon}. \tag{55}$$

Note, that in SUSY QCD, contrary to the case of non-supersymmetric QCD, it coincides with the gauge charge renormalization constant. It is due to supersymmetry, which now protects tree-level coincidence of the $\varepsilon$-scalar and vector coupling constants.

The one-loop nonminimal mass counterterm for $\varepsilon$-scalars reads

$$\delta m_s^2 = -\frac{\alpha_s}{4\pi}\left(\sum_{n_f}\left(\frac{m_f^2}{\mu^2}\right)^{-\varepsilon} m_f^2\left[\frac{2}{\varepsilon}+2\right] - \sum_{n_{\tilde{f}}}\left(\frac{m_{\tilde{f}}^2}{\mu^2}\right)^{-\varepsilon} m_{\tilde{f}}^2\left[\frac{1}{\varepsilon}+1\right] + \left(\frac{m_{\tilde{g}}^2}{\mu^2}\right)^{-\varepsilon} m_{\tilde{g}}^2\left[\frac{2}{\varepsilon}+2\right]\right). \tag{56}$$

It is necessary to insure that the pole mass of $\varepsilon$-scalar is zero. Here $n_f$ denotes sum over different quark flavours and $n_{\tilde{f}}$ used to denote sum over different squarks. It is just this counterterm, that cancels $m_{\mathrm{hard}}^2$ terms occurring in two-loop diagrams and ensures decoupling of large scale physics for physical quantity like quark pole mass. A way to renormalize $\varepsilon$-scalar mass nonminimally was first proposed in [7] and shown to be a consistent renormalization procedure. In fact, there are two theoretically admissible prescriptions to deal with $\varepsilon$-scalar mass renormalization. One is to renormalize $\varepsilon$-scalar mass minimally [20] and another one is the total subtraction of loop corrections to the $\varepsilon$-scalar mass [7]. In the first approach $\varepsilon$-scalar mass becomes a running quantity and physical quantities, like pole mass in our case, will in general depend on it together with other physical renormalized parameters. However, $\varepsilon$-scalar mass is not a physical quantity and it is more correct to treat it as an artifact of our regularization and renormalization schemes. From this point of view the second approach is more natural, if we are going to follow as close as possible the original idea of dimensional reduction suggesting the zero tree-level mass for both $\varepsilon$-scalars and gauge bosons. Of course all renormalization schemes are equivalent and could be related via final recalculation of parameters. An introduction of nonminimal subtraction for $\varepsilon$-scalar mass leads to the fact that bare quark mass is not anymore a renormalization group invariant. The point is that the physical quark mass depends on two mutually correlated bare masses: quark mass itself and $\varepsilon$-scalar mass. Both these quantities are $\mu$ dependent and only the physical pole mass is RG invariant. For precisely this reason the RG equations for the bare quark mass should be written with great care.

# 4 Results of calculation

Here we present the results of our calculation. Corrections to the pole masses are parametrized as follows

$$\frac{M_{\text{pole}}}{m} = 1 + \left(\frac{\Delta m}{m}\right), \tag{57}$$

where $\Delta m$ starts from the one-loop order.

Through calculations we reproduced all known results about radiative corrections to the pole masses of bottom and top quarks: one-loop MSSM [12, 13, 14, 15, 16] and two loop QCD [8, 9] (in $\overline{\text{MS}}$ ) and [7] (in $\overline{\text{DR}}$ ) corrections. As an additional check the renormalization group analysis can be applied to confirm independently the $1/\varepsilon$ and $1/\varepsilon^2$ terms. We get complete agreement between our pole terms and those from RG for $\overline{\text{DR}}$ quark mass renormalization constant. And the last but not least: calculation were performed in the linear gauge with a free QCD gauge parameter. The correction to the pole mass (57) is however appears to be gauge independent.

## 4.1 one-loop result

The corrections to the bottom and top quark pole masses at one loop are well known for a long time for QCD as well as for MSSM. Here we reproduce them only up to terms of order $\mathcal{O}(m^2_{\text{soft}}/m^2_{\text{hard}})$ in large mass expansion procedure. For pure QCD we have:

$$\left(\frac{\Delta m_b}{m_b}\right)^{\text{QCD}} = C_F \frac{\alpha_s}{4\pi}\left[5 - 3\ln\left(\frac{m_b^2}{\mu^2}\right)\right] \tag{58}$$

and squark-gluino MSSM contribution is given by

$$\begin{aligned}
\left(\frac{\Delta m_b}{m_b}\right)^{\tilde{b}\tilde{g}} &= C_F \frac{\alpha_s}{8\pi}\left[-3 + \frac{m^2_{\tilde{b}_1}}{m^2_{\tilde{b}_1} - m^2_{\tilde{g}}} + \frac{m^2_{\tilde{b}_2}}{m^2_{\tilde{b}_2} - m^2_{\tilde{g}}} + 2\ln\left(\frac{m^2_{\tilde{g}}}{\mu^2}\right)\right. \\
&+ \frac{m^2_{\tilde{b}_1}}{m^2_{\tilde{b}_1} - m^2_{\tilde{g}}}\left(2 - \frac{m^2_{\tilde{b}_1}}{m^2_{\tilde{b}_1} - m^2_{\tilde{g}}} - 2\sin(2\theta_{\tilde{b}})\frac{m_{\tilde{g}}}{m_b}\right)\ln\left(\frac{m^2_{\tilde{b}_1}}{m^2_{\tilde{g}}}\right) \\
&+ \left.\frac{m^2_{\tilde{b}_2}}{m^2_{\tilde{b}_2} - m^2_{\tilde{g}}}\left(2 - \frac{m^2_{\tilde{b}_2}}{m^2_{\tilde{b}_2} - m^2_{\tilde{g}}} + 2\sin(2\theta_{\tilde{b}})\frac{m_{\tilde{g}}}{m_b}\right)\ln\left(\frac{m^2_{\tilde{b}_2}}{m^2_{\tilde{g}}}\right)\right]
\end{aligned} \tag{59}$$

which coincide with the results of [12, 13, 14, 15, 16]. To obtain similar expressions for the case of top one should perform obvious substitution $b \to t$. Results for squark-chargino and higgs-quarks loops also known, but we interesting in this paper only in $\mathcal{O}(\alpha_s)$ and $\mathcal{O}(\alpha_s^2)$ correction.

## 4.2 Two-loop result

The two-loop QCD result in $\overline{\mathrm{DR}}$-scheme for $b$-quark (with four light quarks)

$$\begin{aligned}
\left(\frac{\Delta m_b}{m_b}\right)^{\mathrm{QCD}} &= C_F \left(\frac{\alpha_s}{4\pi}\right)^2 \Bigg\{ -\frac{623}{18} - 8\,\zeta_2 + \ln^2\left(\frac{m_b^2}{m_t^2}\right) - \frac{13}{3}\ln\left(\frac{m_b^2}{m_t^2}\right) \\
&+C_F\left[-\frac{59}{8} + 30\,\zeta_2 - 48\,\ln(2)\,\zeta_2 + 12\,\zeta_3 + \frac{3}{2}\ln\left(\frac{m_b^2}{\mu^2}\right) + \frac{9}{2}\ln^2\left(\frac{m_b^2}{\mu^2}\right)\right] \\
&+C_A\left[\frac{1093}{24} - 8\,\zeta_2 + 24\,\ln(2)\,\zeta_2 - 6\,\zeta_3 - \frac{179}{6}\ln\left(\frac{m_b^2}{\mu^2}\right) + \frac{11}{2}\ln^2\left(\frac{m_b^2}{\mu^2}\right)\right] \\
&+ 26\,\ln\left(\frac{m_b^2}{\mu^2}\right) - 6\ln^2\left(\frac{m_b^2}{\mu^2}\right)\Bigg\}
\end{aligned} \tag{60}$$

and for $t$-quark (with five light quarks)

$$\begin{aligned}
\left(\frac{\Delta m_t}{m_t}\right)^{\mathrm{QCD}} &= C_F \left(\frac{\alpha_s}{4\pi}\right)^2 \Bigg\{ -43 - 12\,\zeta_2 + 26\,\ln\left(\frac{m_t^2}{\mu^2}\right) - 6\ln^2\left(\frac{m_t^2}{\mu^2}\right) \\
&+C_F\left[-\frac{59}{8} + 30\,\zeta_2 - 48\,\ln(2)\,\zeta_2 + 12\,\zeta_3 + \frac{3}{2}\ln\left(\frac{m_t^2}{\mu^2}\right) + \frac{9}{2}\ln^2\left(\frac{m_t^2}{\mu^2}\right)\right] \\
&+ C_A\left[\frac{1093}{24} - 8\,\zeta_2 + 24\,\ln(2)\,\zeta_2 - 6\,\zeta_3 - \frac{179}{6}\ln\left(\frac{m_t^2}{\mu^2}\right) + \frac{11}{2}\ln^2\left(\frac{m_t^2}{\mu^2}\right)\right]\Bigg\}
\end{aligned} \tag{61}$$

coincide with the first terms of the expansion in $m_b/m_t$ from [7]

We find very simple answer for one particular limit. Namely, in the case, when $m_{\tilde{t}_1} = m_{\tilde{t}_2} = m_{\tilde{b}_1} = m_{\tilde{b}_2} = m_{\tilde{g}} = m_{\tilde{u}_{1,2}} = m_{\tilde{d}_{1,2}} = m_{\mathrm{SUSY}} = M$, the expression looks like

$$\begin{aligned}
\left(\frac{\Delta m_q}{m_q}\right)^{\mathrm{MSSM}} &= \left(\frac{\alpha_s}{4\pi}\right)^2 \Bigg\{ C_F^2\left[\frac{23}{24} - \frac{13}{6}\ln\left(\frac{M^2}{\mu^2}\right) + \frac{1}{2}\ln^2\left(\frac{M^2}{\mu^2}\right) - 3\ln\left(\frac{M^2}{\mu^2}\right)\ln\left(\frac{m_q^2}{\mu^2}\right)\right] \\
&+ C_F\left[\frac{47}{3} + 20\ln\left(\frac{M^2}{\mu^2}\right) + 6\ln^2\left(\frac{M^2}{\mu^2}\right) - 6\ln\left(\frac{M^2}{\mu^2}\right)\ln\left(\frac{m_q^2}{\mu^2}\right)\right] \\
&+ C_F C_A\left[\frac{175}{72} + \frac{41}{6}\ln\left(\frac{M^2}{\mu^2}\right) - \frac{1}{2}\ln^2\left(\frac{M^2}{\mu^2}\right) - 2\ln\left(\frac{M^2}{\mu^2}\right)\ln\left(\frac{m_q^2}{\mu^2}\right)\right]\Bigg\}
\end{aligned} \tag{62}$$

Since the complete formulae[8] are too large to be presented here, we present only the influence of our results on the masses of $b$- and $t$-quark and the spectra of supersymmetric particles, obtained as solutions of RG equations. We incorporated our formulae into SoftSUSY code [6] and analysed the difference in the predictions for particle masses obtained with and without our corrections to $b$- and $t$-quark pole masses.

$$\Delta_p = \frac{m_p^{\mathrm{CODE}} - m_p^{\mathrm{CODE+Corrections}}}{m_p^{\mathrm{CODE}}}, \tag{63}$$

[8] They could be found in source archive file for this paper, submitted to HEP database

where $p$ stands for quarks or different supersymmetric particles.

In Fig.6 we present the variation of $\Delta_p$ as a function of $\tan\beta$ for $m_0 = 400\,\text{GeV}$, $m_{1/2} = 400\,\text{GeV}$, $A_0 = 0$ and $M_{GUT} = 1.9 \times 10^{16}\,\text{GeV}$. We plotted $\Delta_p$ for five different cases:

• **Bottom 1L** - difference between results obtained with code including and excluding MSSM one loop corrections to the pole masses of $b$-quarks

• **Top 1L** - the same as above for $t$-quarks

• **Bottom** - difference between results obtained with code containing two loop MSSM $\mathcal{O}(\alpha_s^2)$ correction to $b$-quark together with MSSM one loop corrections to the pole masses of $b$- and $t$-quarks and results obtained with code containing MSSM one loop corrections to the pole masses of b- and $t$-quarks

• **Top** - the same as above for $t$-quarks

• **All** - difference between results obtained with code containing two loop MSSM $\mathcal{O}(\alpha_s^2)$ correction together with MSSM one loop corrections to the pole masses of b- and $t$-quarks and results obtained with code without any MSSM (one loop and two loop) corrections to the pole masses of b- and $t$-quarks.

From this figure one may see, that two loop correction for the relation between $b$-quark pole and $\overline{\text{DR}}$ masses gives sizeable contribution only to the solution obtained with SoftSUSY program for $\overline{\text{DR}}$ $b$-quark mass at $M_Z$ and almost do not influence spectra of SUSY particles. However, if one consider large $\tan\beta$ it becomes impotent for other particle spectra and thus should be taken into account in accurate RG analysis. Note, that our two-loop SUSY QCD correction for a wide range of parameter space is always less, then similar 1-loop contribution. The two-loop MSSM correction for $t$-quark pole mass, on the other hand, contributes to almost all SUSY particle masses, obtained as a solution of renormalization group equations with universal boundary conditions at GUT scale. This 2-loop correction greater then 1-loop MSSM contribution almost everywhere in MSSM parameter space. This could be attributed to potentially large corrections coming from higher terms in $m_t/M_{SUSY}$ expansion.

To see the dependence of $\Delta_p$ on $m_0$, $m_{1/2}$ and $A_0$ we use a set of benchmark points and parameter lines in the MSSM parameter space from [26, 27, 28, 29] which corresponds to different scenarios in the search for Supersymmetry at present and future colliders. While a detailed scan over the more-than-hundred-dimensional parameter space of the MSSM is clearly not practicable, even a sampling of the three- (four-)dimensional parameter space of $m_0$, $m_{1/2}$ and $A_0$ ($\tan\beta$) is beyond the present capabilities for phenomenological studies, especially when one tries to simulate experimental signatures of supersymmetric particles within detectors. For this reason one often resorts to specific benchmark scenarios, i.e. one studies only specific parameter points or at best samples a one-dimensional parameter space (the latter is sometimes called a model line [28]), which exhibit specific characteristics of the MSSM parameter space.

Some recent proposals for SUSY benchmark scenarios may be found in [26, 27, 29]. We refer to the "Snowmass Points and Slopes" (SPS) [28, 29], which consist of model lines ("slopes"), i.e. continuous sets of parameters depending on one dimensionful parameter and specific benchmark points, where each model line goes through one of the benchmark points.

In Fig.7 we present the results for $\Delta_p$ as a function defined along **Model Line A**

$$m_0 = -A_0 = 0.4\, m_{1/2}, \quad m_{1/2} \text{ varies}, \quad \tan\beta = 10, \quad \mu > 0.$$

where benchmark point is

$$m_0 = 100\,\text{GeV}, \quad m_{1/2} = 250\,\text{GeV}, \quad A_0 = -100\,\text{GeV}, \quad \tan\beta = 10, \quad \mu > 0.$$

Here and below $\mu$ is the supersymetric Higgs mass parameter

As one can easily see from Fig.7, increasing $m_{1/2}$ at fixed $\tan\beta$ 2-loop SUSY QCD correction to $b$-quark pole mass almost do not contribute to heavy supersymmetric particle spectra. On the other hand, our correction to $t$-quark pole mass is more important at small values of $m_{1/2}$, where it becomes even larger then corresponding 1-loop contribution.

In Fig.8 we present the results for $\Delta_p$ defined as a function along **Model Line C**

$$m_0 = m_{1/2}, \quad m_{1/2} \text{ varies}, \quad A_0 = 0, \quad \tan\beta = 35, \quad \mu > 0.$$

where benchmark point is

$$m_0 = 300\,\mathrm{GeV}, \quad m_{1/2} = 300\,\mathrm{GeV}, \quad A_0 = 0\,\mathrm{GeV}, \quad \tan\beta = 35, \quad \mu > 0.$$

Situation with Fig.8 is analogous to that of Fig.7, even if we are increasing $\tan\beta$ from 10 up to 35. An importance of our corrections increases only at very large values of $\tan\beta$ (50 and more), as one can see from Fig.6 .

For such large $\tan\beta$ we just used some particular values $m_0$ and $m_{1/2}$ from allowed region from [30]. We use next parametrization

$$m_0 = m_{1/2}, \quad m_{1/2} \text{ varies}, \quad A_0 = 0, \quad \tan\beta = 55, \quad \mu > 0.$$

The results for $\Delta_p$ are shown in Fig.9 and we see, that in this case our 2-loop correction to $b$-quark pole mass gives sizeable contribution to the masses of gluino and sbottoms.

# 5 Conclusion

We presented results for two-loop MSSM corrections to the relation between pole and running masses of heavy quarks up to $\mathcal{O}(\alpha_s^2)$ order. We provided a detail analysis of the value of these corrections in different regions of parameter space and discussed their impact on SUSY particle spectra. We would like to note, that our results presented here may be also included in the codes of programs like Isajet[31], SuSpect[32] and used for predictive phenomenological analyses like [30], [33], [34], [35]. In one of our next papers we suppose to calculate missing corrections from stop-chargino loops to $b$-quark pole mass as well as supply more terms in the expansion in a relation between top pole and $\overline{\mathrm{DR}}$ masses.

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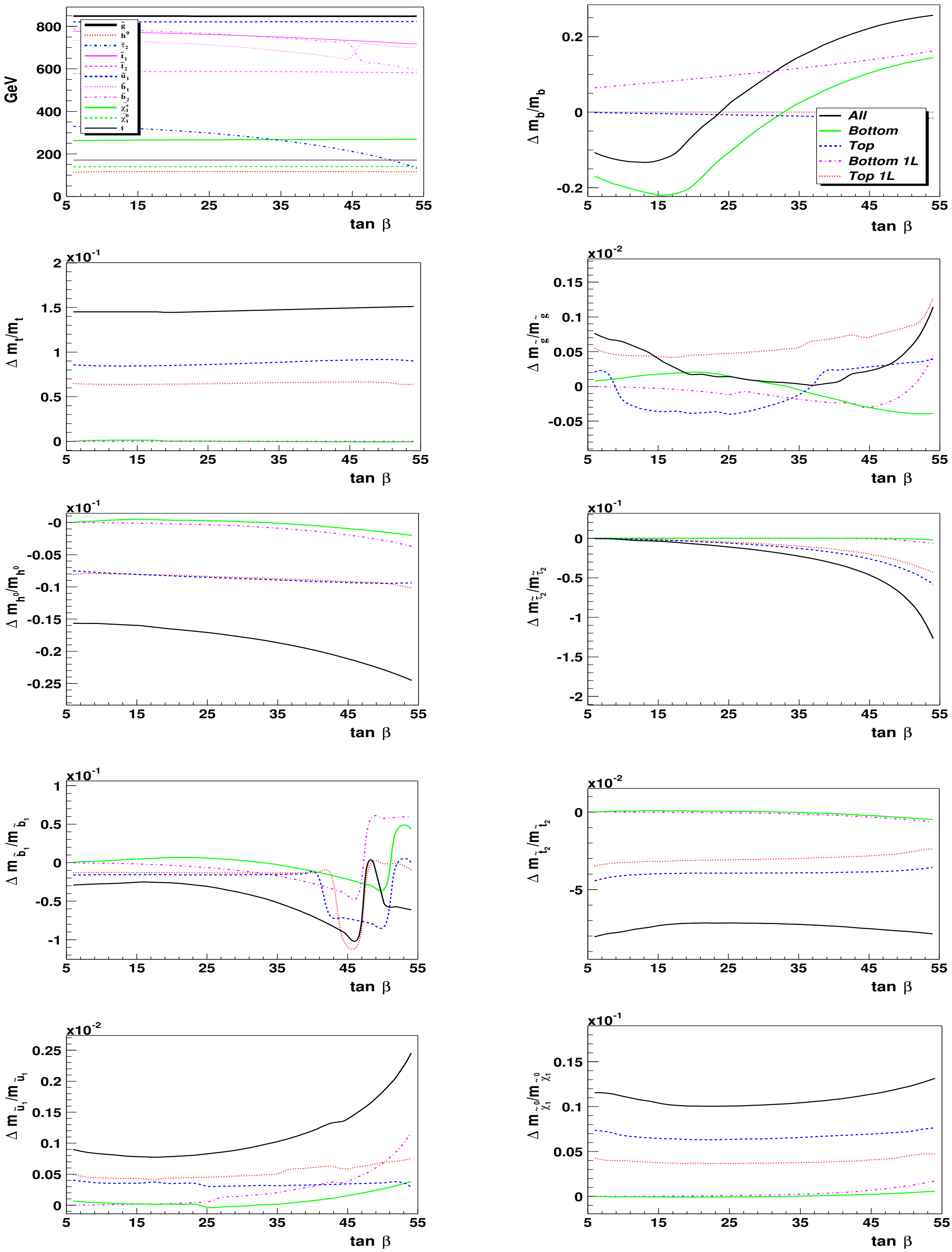


Figure 6: $\Delta_p$ as a function of $\tan\beta$. In the first picture we plotted a specter of masses

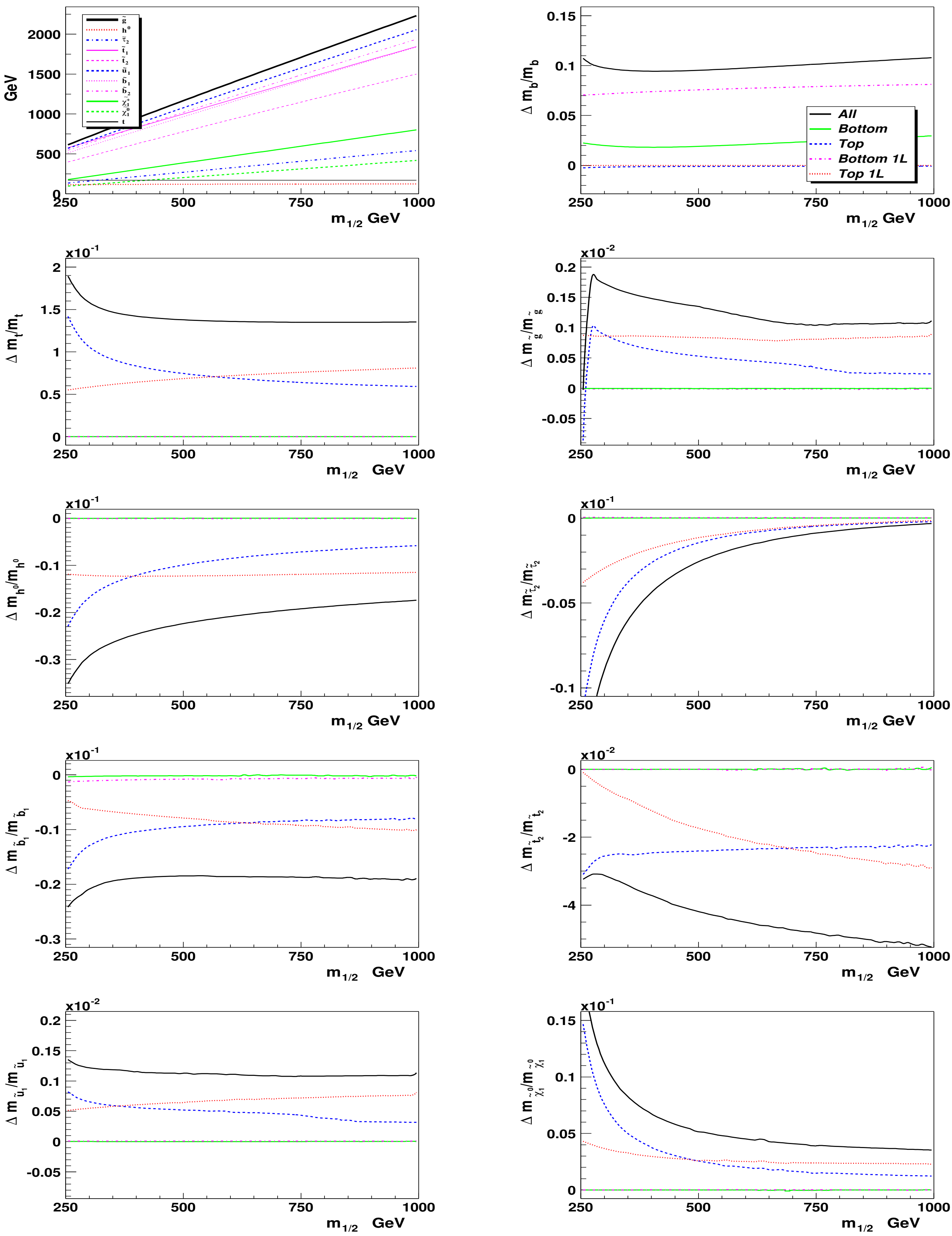


Figure 7: $\Delta_p$ as a function defined along **Model Line A** from corrections on $m_{1/2}$. In the first picture we plotted a specter of masses

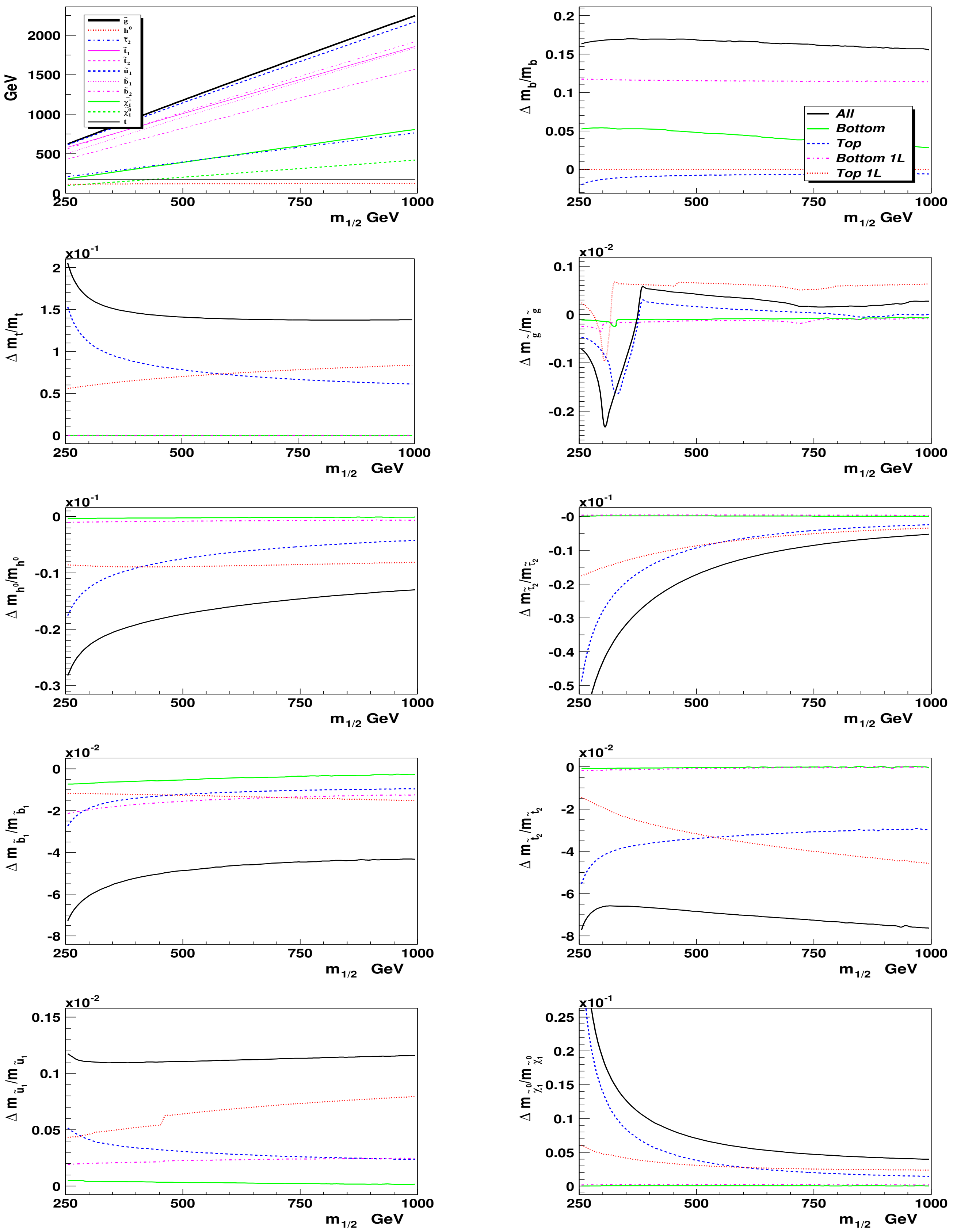


Figure 8: $\Delta_p$ as a function defined along **Model Line C** from corrections on $m_{1/2}$. In the first picture we plotted a specter of masses

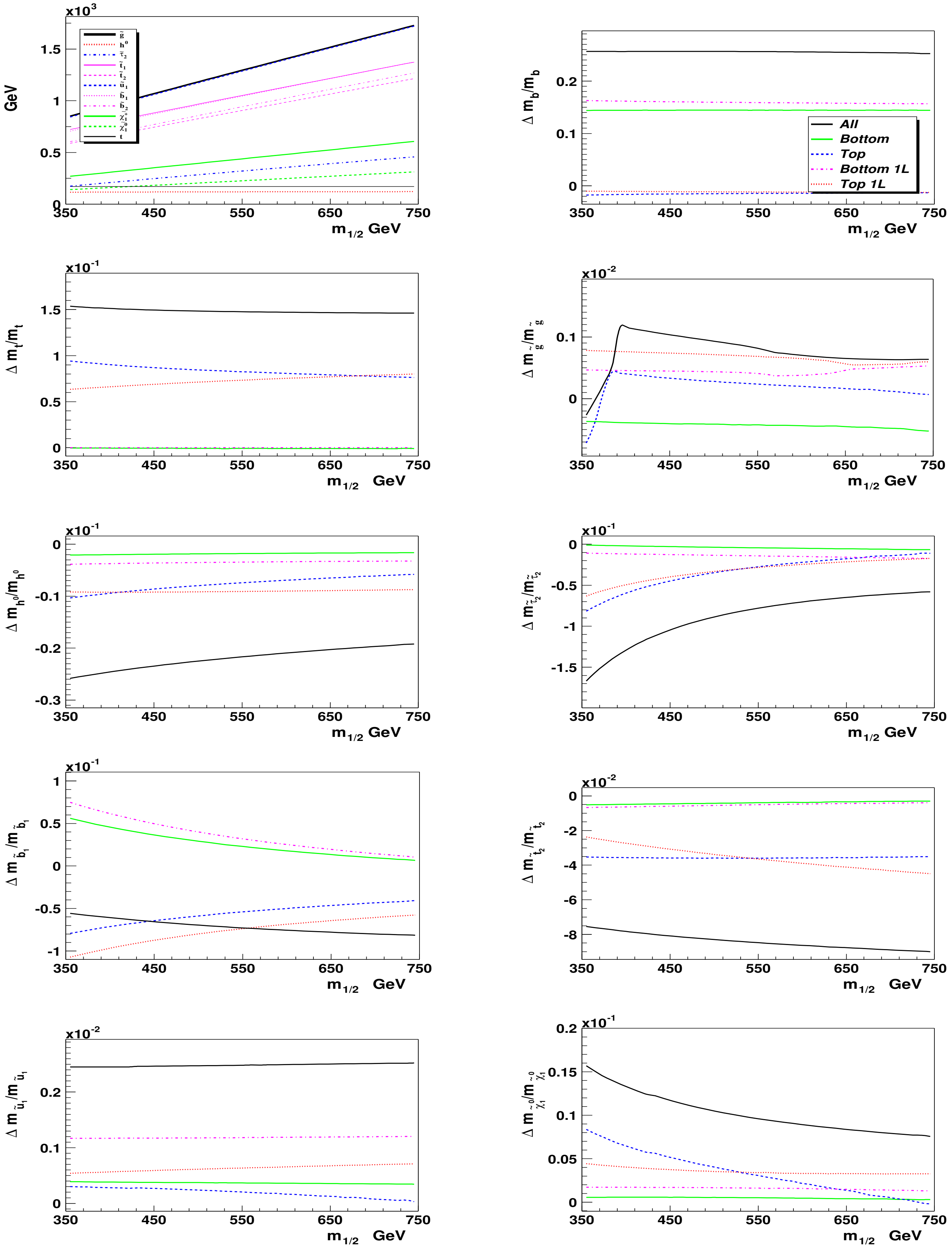


Figure 9: $\Delta_p$ at large value of $\tan\beta = 55$ as a function of $m_{1/2}$. In the first picture we plotted a specter of masses