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New LEP bounds on *B*-violating scalar couplings: *R*-parity violating supersymmetry or diquarks

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ABSTRACT

We use the precision electroweak data at LEP to place bounds on B-violating Yukawa couplings, two theoretically appealing examples being provided by R-parity–violating supersymmetry and diquarks. The couplings involving the third generation quarks are most severely constrained. These bounds are complementary to those obtained from low-energy processes.

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One of the goals of the precision measurements on the Z peak at LEP is to probe scenarios beyond the Standard Model (SM) predicting the existence of new particles coupling to the Z. The experimental sensitivity is enhanced if these particles have tree-level couplings to the SM fermions as well. Particular examples are provided by R-violating supersymmetry, as also by diquarks, the existence of which can be motivated in the context of both Grand Unified Theories as well as composite models[1] [1]. Currently, there exist phenomenological bounds — from low-energy data — on only some of these couplings. We aim to use the LEP data on Z partial widths to impose complementary bounds.

A diquark can, in general, be defined as an elementary integral-spin particle with a baryon number $|B| = 2/3$ and lepton number $|L| = 0$ and coupling to a pair of quarks. Assuming the SM fermion content, the diquarks may transform as $\mathbf{3}$ or $\bar{\mathbf{6}}$ under $SU(3)_C$; as triplet or singlet under $SU(2)_L$; and can have electric charges $|Q_D| = 1/3, 2/3$ or $4/3$. In this letter, we restrict ourselves to a discussion of scalar diquarks only. The Yukawa interaction terms can then be parametrized as:

$$\mathcal{L}_Y = h_{ij}^{(A)} \overline{q_i^c} P_{L,R} q_j \phi_A + \text{h.c.} \; , \tag{1}$$

where i and j denote the flavour generation indices, A labels the diquark-type and $h_{ij}^{(A)}$ are the corresponding Yukawa couplings. The relevant projection operator $P_{L,R}$ depends upon the gauge transformation properties of the concerned fields. All such possible contributions within the texture of the SM are listed in Table 1. The coupling of the diquark to the gauge bosons is obviously determined by the quantum numbers. In addition, it will have self-coupling terms, and also couplings with the SM Higgs, but these are of no consequence to us.

The above Yukawa interaction is of a rather general type, parts of which are mimiced by the B-violating parts of the R-parity-violating ($\not{R}_p$) Yukawa interaction in supersymmetric theories [2, 3]. Representable as $R = (-1)^{3B+L+2S}$, where B, L, S are the baryon number, lepton number and the intrinsic spin of the field, respectively, R has a value of $+1$ for all SM particles and -1 for all their superpartners. The B-violating part of the $\not{R}_p$ superpotential is

$$\mathcal{W} = \lambda''_{ijk} U_i^c D_j^c D_k^c \; , \tag{2}$$

where U_i^c and D_j^c are the up- and down-type quark superfields respectively. The couplings λ''_{ijk} are obviously antisymmetric in the last two indices. The above interaction can be rewritten in terms of the component fields as:

$$\mathcal{L}_{\not{R}_p} = \lambda''_{ijk} \left(u_i^c d_j^c \tilde{d}_k^* + u_i^c \tilde{d}_j^* d_k^c + \tilde{u}_i^* d_j^c d_k^c \right) + \text{h.c.} \tag{3}$$

It is quite apparent that eq.(3) replicates parts of eq.(1) when the quark fields in the latter are $SU(2)_L$-singlets. To be precise, we have a situation where scalars of both type ϕ_8 ($\tilde{d}_i$) and ϕ_{10} ($\tilde{u}_i$) exist simultaneously.

[1] It may be noted that these theories predict the existence of leptoquarks as well. However, the presence of certain types of leptoquarks in association with diquarks would lead to rapid proton decay. In this letter we assume that leptoquark couplings, if existing, are highly suppressed.

Scalar type	Coupling	$SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ transformation	Remarks
ϕ_1	$\overline{(Q_{Li})^c} Q_{Lj} \phi_1$	$(\bar{6}, 3, -1/3)$	Generation symmetric
ϕ_2	$\overline{(Q_{Li})^c} Q_{Lj} \phi_2$	$(3, 3, -1/3)$	Generation antisymmetric
ϕ_3	$\overline{(Q_{Li})^c} Q_{Lj} \phi_3$	$(\bar{6}, 1, -1/3)$	Generation antisymmetric
ϕ_4	$\overline{(Q_{Li})^c} Q_{Lj} \phi_4$	$(3, 1, -1/3)$	Generation symmetric
ϕ_5	$\overline{(u_{Ri})^c} u_{Rj} \phi_5$	$(\bar{6}, 1, -4/3)$	Generation symmetric
ϕ_6	$\overline{(u_{Ri})^c} u_{Rj} \phi_6$	$(3, 1, -4/3)$	Generation antisymmetric
ϕ_7	$\overline{(u_{Ri})^c} d_{Rj} \phi_7$	$(\bar{6}, 1, -1/3)$	No symmetry property
ϕ_8	$\overline{(u_{Ri})^c} d_{Rj} \phi_8$	$(3, 1, -1/3)$	No symmetry property
ϕ_9	$\overline{(d_{Ri})^c} d_{Rj} \phi_9$	$(\bar{6}, 1, 2/3)$	Generation symmetric
ϕ_{10}	$\overline{(d_{Ri})^c} d_{Rj} \phi_{10}$	$(3, 1, 2/3)$	Generation antisymmetric

Table 1: The possible diquark couplings within the SM quark content.

It may be noted that there are important cosmological constraints [4] on the R-parity–violating Yukawa interactions. Requiring that GUT-scale baryogenesis does not get washed out imposes $\lambda'' \ll 10^{-7}$ generically, though these bounds are model dependent and can be evaded [5]. Evidently, similar considerations also hold for diquarks.

As for phenomenological bounds on the diquarks, very little exists. The preferred arena for their production is obviously a hadron collider. However, the large QCD background to the signal for diquark pair-production results in poor detectability. One might think that the situation is slightly better for $\not{R}_p$ interactions, as there are other channels through which the squarks might decay[2]. It must be remembered though, that in the presence of $\not{R}_p$ couplings, the lightest supersymmetric partner is no longer stable, and many of the bounds do not hold. Moreover, the detection of a squark *per se* is not a valid indicator of a $\Delta B = 1$ $\not{R}_p$ interaction.

Some bounds have been obtained from low-energy processes though. For example, the data on neutral meson mixing or that on CP–violation in the K–sector can tightly constrain certain products of such couplings [3]. As for individual bounds, strong ones exist only for certain cases where both quarks are light. These are derived from the non-observance of neutron–antineutron oscillations [7, 8] and double nucleon decay into kaons of identical strangeness [8].

In this paper, we consider the effects of the Yukawa interactions of the diquarks (and

[2]Indeed, the CDF experiment at the Tevatron has quoted lower bounds on the squark mass of ~ 150 GeV, modulo certain assumptions about the supersymmetric parameter space [6].

also of the squarks through their B-violating Yukawa interactions) on the process $Z \to q\bar{q}$, where q is a quark. These interactions proceed through triangle and self-energy diagrams with Z, q and $\bar{q}$ as external legs. In each triangle and self-energy diagram the internal lines are the diquark(s) and the quark(s) which couple(s) to the diquark. A typical set of triangle and self-energy diagrams with a generic diquark (quark) ϕ (Q) are shown in Fig. 1. The magnitude of the contribution grows as the mass of the internal quark in this case. Hence we focus our attention on the terms involving the top quark.

The tree-level Z couplings to the left- and right-handed fermions can be parametrized as

$$M_\mu^{\rm tree} = e\bar{q}(p')\gamma_\mu(a_L^q P_L + a_R^q P_R)q(p). \tag{4}$$

where

$$\begin{aligned} a_L^q &= (t_3^q - Q_q s_W^2)/s_W c_W, \\ a_R^q &= -Q_q s_W/c_W. \end{aligned} \tag{5}$$

The Z couplings to the charge-conjugated fermions (q^c) are, therefore, given by

$$a_L^{q^c} = -a_R^q, \qquad\qquad a_R^{q^c} = -a_L^q. \tag{6}$$

We compute the self-energy and the vertex correction diagrams in terms of the Passarino-Veltman B- and C-functions [9], corresponding to the two- and three-point integrals[3]. Assuming for the sake of presentation that in Fig. 1 the external fermions lines are *left-handed*, it is easy to see that only the couplings of the left-handed fermions to the Z are modified, so that the new contribution to the amplitude is given by

$$M_\mu^{(i)} = \frac{eh^2 N_c}{16\pi^2}\, \bar{q}(p')\, \gamma_\mu\, A_i\, P_L\, q(p), \tag{7}$$

where $i = 1, 2, 3$, h is a generic Yukawa coupling, and N_c is the colour factor which is 3 (2) when ϕ belongs to a $\bar{\mathbf{6}}$ ($\mathbf{3}$) of $SU(3)_C$. The A_i's are given by,

$$\begin{aligned} A_1 &= a_L^Q m_t^2 C_0 - a_R^Q\{m_Z^2(C_{22} - C_{23}) + (d-2)C_{24}\}, \\ A_2 &= -2\frac{t_3^\phi - Q_\phi s_W^2}{s_W c_W}\tilde{C}_{24}, \\ A_3 &= a_L^q B_1. \end{aligned} \tag{8}$$

Here $A_{1,2}$ denote the contributions from the first and the second triangle diagrams respectively, and the contributions of the two self-energy diagrams are jointly denoted by A_3. In A_2, we use $\tilde{C}_{24}$ to distinguish it from the C_{24} appearing in A_1, as the structures of the propagators for the two triangle diagrams are different. In the expression for A_1, d denotes the space-time dimension. Although the individual diagrams are divergent, their sum is finite. The asymptotic forms for A_i can be found in refs. [11], which deal with similar bounds on lepton number–violating Yukawa couplings.

To compare our results with the experimental numbers we use the following observables:

[3] We use the numerical codes developed by Mukhopadhyaya and Raychaudhuri in the context of [10].

1. $R_l = \Gamma_{\rm had}/\Gamma_l$, which is stable under variation of the top-quark mass[4]. Recent measurements [12] give $R_l^{\rm exp} = 20.795 \pm 0.04$ whereas the SM prediction is $R_l^{\rm SM} = 20.786$ for a choice of $m_t = 175$ GeV, $m_H = 300$ GeV and $\alpha_S = 0.12$.

2. $R_b = \Gamma_b/\Gamma_{\rm had}$, which has a quadratic top mass dependence. From ref.[12], $R_b^{\rm exp} = 0.2202 \pm 0.0020$ and $R_b^{\rm SM} = 0.2158$ for the above choice of input parameters.

3. $R_c = \Gamma_c/\Gamma_{\rm had}$, which again has a quadratic top mass dependence. Ref.[12] quotes $R_c^{\rm exp} = 0.1583 \pm 0.0098$ and $R_c^{\rm SM} = 0.172$ for the same choice of input parameters as above.

In Figs. 2–4, we plot respectively the deviations δR_l, δR_b and δR_c caused by the presence of ϕ_A ($A = 2, 4, 6, 8$) as a function of the scalar mass m_ϕ. In each case, we assume that only the corresponding Yukawa coupling is non-zero and equals unity. In Table 2 we show limits on the Yukawa couplings of each scalar type for a common scalar mass of 100 GeV.

We summarize our results below:

- The bounds obtained from R_l are always better than those from R_b, mainly because R_l is measured with an accuracy of 0.2%, while R_b is measured with an accuracy of not better than 1% at this stage. Note, however, that many of the diquark couplings lead to a negative contribution to R_b. Thus, with an improvement in this determination, all such couplings can be constrained to a greater extent.

- We do not use the experimental numbers on R_c to put any bounds, as the experimental errors are relatively large. Notice, though, that *all* of the diquark couplings lead to a *positive* contribution to R_c. Thus, if the experimental value of R_c continues to stay below the SM prediction once the errors are reduced significantly, this measurement will disfavour such scalars.

- Only those contributions have been shown in figures which are induced by $SU(3)_C$-triplet scalars. In cases where a $\bar{\mathbf{6}}$ scalar may couple as well, it is evident that its contribution is similar to that of the $\mathbf{3}$, except for a colour enhancement factor of 3/2. The corresponding bounds on its Yukawa couplings are thus stronger by a factor of $\sqrt{3/2}$, as can be seen from Table 2.

- Scalars which couple to the top quark are constrained more stringently than others because the largest contributions[5] to δR_l, δR_b or δR_c are $\mathcal{O}(m_Q^2/m_\phi^2)$, where Q represents the heavy quark in the diagram. Consequently, the bounds on $h_{ij}^{(9)}$ and $h_{ij}^{(10)}$ obtained by our method are an order of magnitude weaker than the rest; hence we do not include those cases either in the figures or in Table 2.

- Since the tree level predictions for Γ_d and Γ_s are nearly the same and also since $m_d, m_s \ll m_t, m_\phi$, $h_{13}^{(A)} \simeq h_{23}^{(A)}$ for all A.

[4]We assume leptonic universality.

[5]We find that the next-to-leading terms can be important in some cases, though.

Coupling	Bounds from R_l			Bounds from R_b		
	(1σ)	(2σ)	(3σ)	(1σ)	(2σ)	(3σ)
$h^{(1)}_{33}$	0.35	0.53	0.66	-	-	0.41
$h^{(1)}_{13}$, $h^{(1)}_{23}$	0.89	1.19	1.44	-	-	2.39
$h^{(2)}_{13}$, $h^{(2)}_{23}$	1.09	1.46	1.76	-	-	2.93
$h^{(3)}_{13}$, $h^{(3)}_{23}$	0.54	0.82	1.03	1.89	2.17	2.41
$h^{(4)}_{33}$	0.60	0.91	1.14	-	-	0.70
$h^{(4)}_{13}$, $h^{(4)}_{23}$	0.66	1.00	1.26	2.32	2.65	2.95
$h^{(5)}_{13}$, $h^{(5)}_{23}$	0.98	1.31	1.58	-	-	1.73
$h^{(6)}_{13}$, $h^{(6)}_{23}$	1.19	1.61	1.94	-	-	2.12
$h^{(7)}_{33}$	1.12	1.69	2.12	-	-	1.31
$h^{(7)}_{31}$, $h^{(7)}_{32}$	1.11	1.68	2.09	4.94	5.66	6.29
$h^{(8)}_{33}$	1.37	2.08	2.60	-	-	1.60
$h^{(8)}_{31}$, $h^{(8)}_{32}$	1.36	2.05	2.57	6.05	6.93	7.71
λ''_{313}, λ''_{323}	0.97	1.46	1.83	-	-	1.89
λ''_{312}	0.96	1.45	1.82	4.28	4.90	5.45

Table 2: The upper bounds on the diquark (and $\not{R}_p$) couplings for $m_\phi = 100$ GeV. Since $R_b^{\rm SM}$ is itself inconsistent with $R_b^{\rm exp}$ at the 2σ level, for most of the cases, R_b does not give a bound at the 1σ or 2σ level.

- For $A = 1$, h_{33} is more strongly constrained than h_{13} or h_{23}. The reasons are twofold. For one, the presence of the top in both quark doublets serves to enhance the effect for h_{33}. Furthermore, in the case of h_{13} (and similarly for h_{23}), while $\delta\Gamma_d < 0$, its effect is negated to an extent by the fact that $\delta\Gamma_u > 0$. This cancellation is obviously absent for h_{33}.

- As is apparent from eq.(3), the $\not{R}_p$ effects can be realized by a linear combination of the contributions due to scalars of the type ϕ_8 and ϕ_{10}. For $i = 3$ the dominance of the former type is overwhelming due to the presence of the heavy top quark in the loop[6]. In Table 2, we show the bounds on λ''-couplings separately although we do not include them in figures. That the bounds are a factor of $\approx \sqrt{2}$ stronger than

[6]For the same reason, we may safely ignore the $\tilde{t}_L$–$\tilde{t}_R$ mixing in our computations.

those on $h_{3i}^{(8)}$ can be understood from the fact that $R\!\!\!/_p$couplings lead to a change in the Z partial width into two quark channels, unlike only one for the diquark, as is evident from a comparison of eq.(1) and eq.(3).

- The bounds on λ''-couplings obtained [7, 8] from the *assumption* of perturbative unification are ~ 1.25, which are at par with the phenomenological bounds we obtain.

In conclusion, we have obtained phenomenological bounds on B-violating Yukawa couplings from LEP data on the Z partial widths. λ''-type Yukawa couplings in $R\!\!\!/_p$ supersymmetric theories and the Yukawa couplings of the diquarks with the standard quarks are two viable candidates of the above type. Our method leads to significant bounds on those couplings which involve the third generation quarks and hence are complementary to those obtained from n–$\bar{n}$ oscillation or other low-energy processes. The limits we obtain will improve with the accumulation of more data at LEP.

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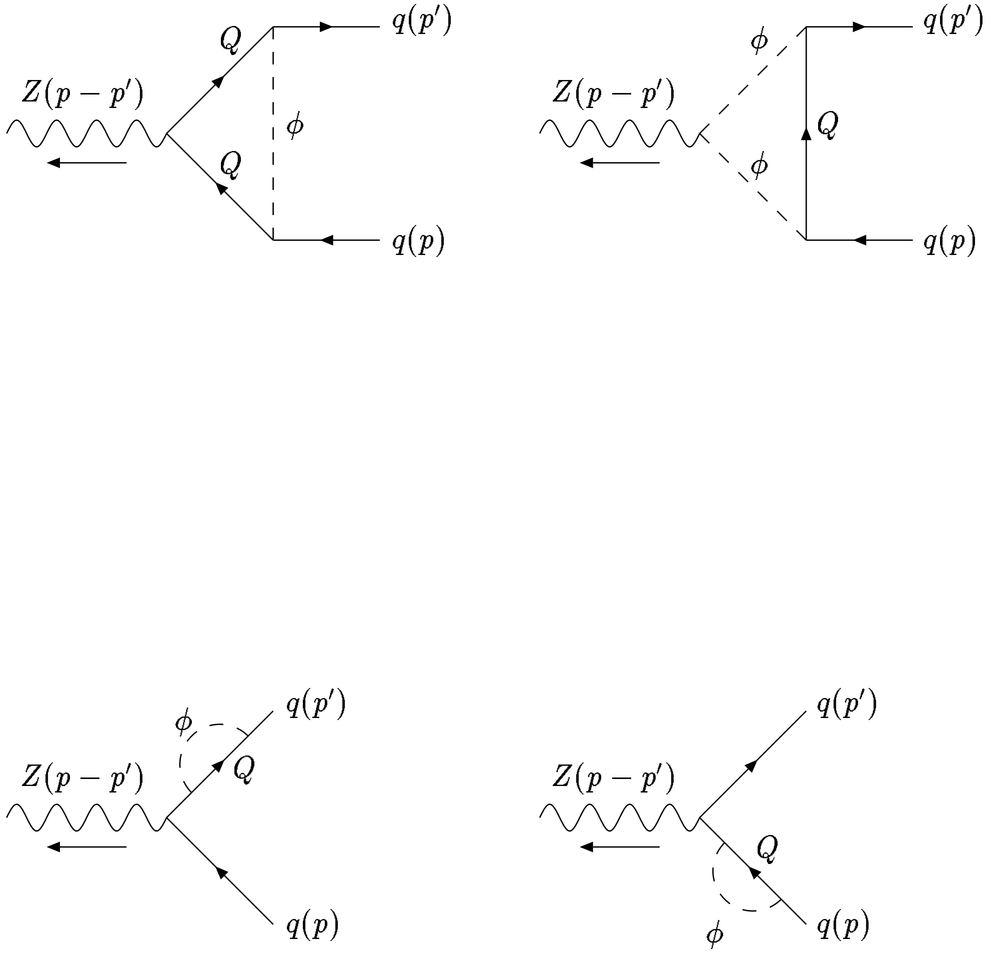

Figure 1: Typical diquark–induced diagrams leading to additional contribution to $\Gamma(Z \to q\bar{q})$.

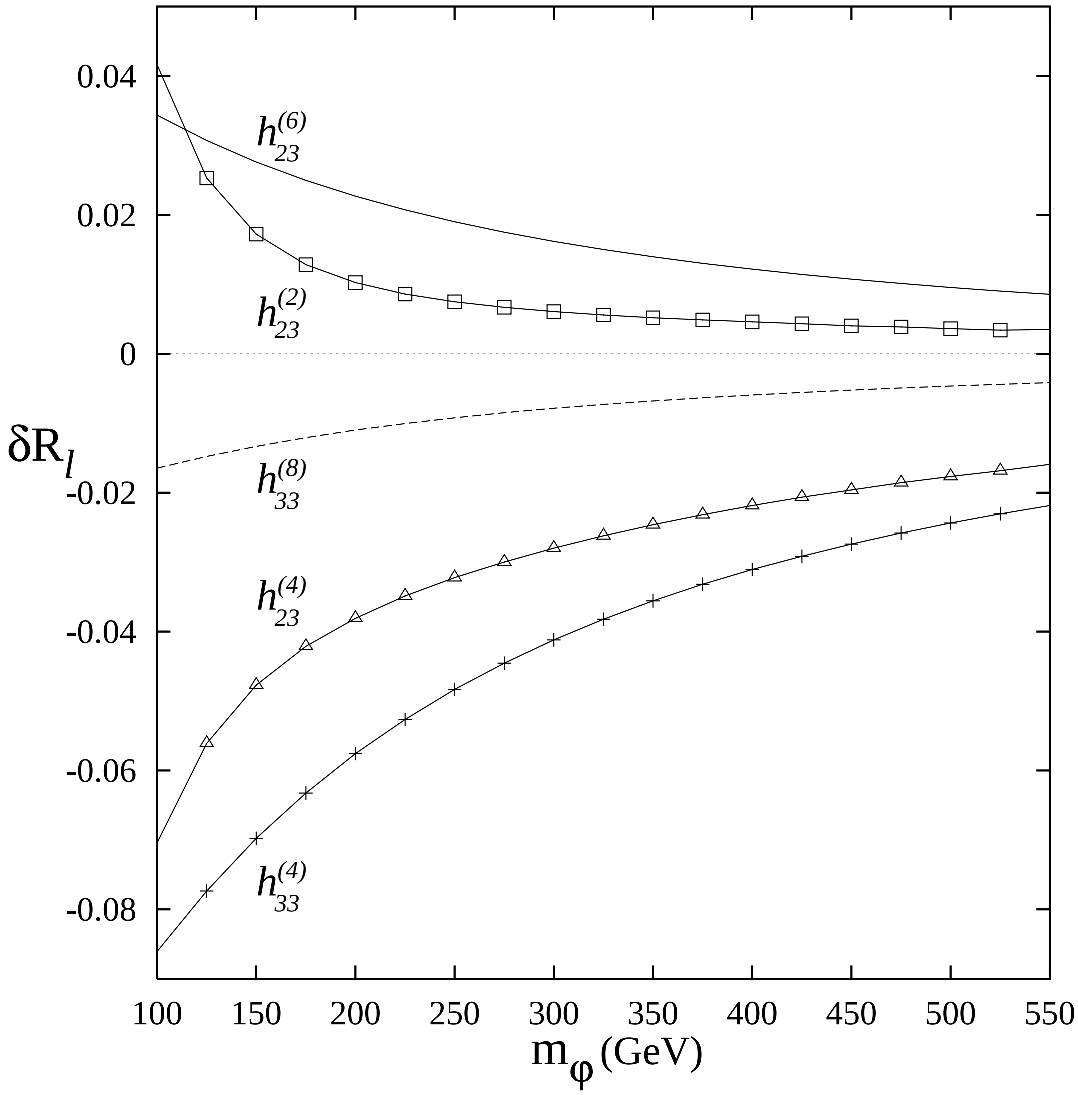

Figure 2: The contribution to R_l in the presence of different B-violating couplings with $SU(3)_c$-triplet scalars. For each individual curve, the concerned Yukawa coupling has been assumed to be unity while all other B-violating couplings are held to be zero.

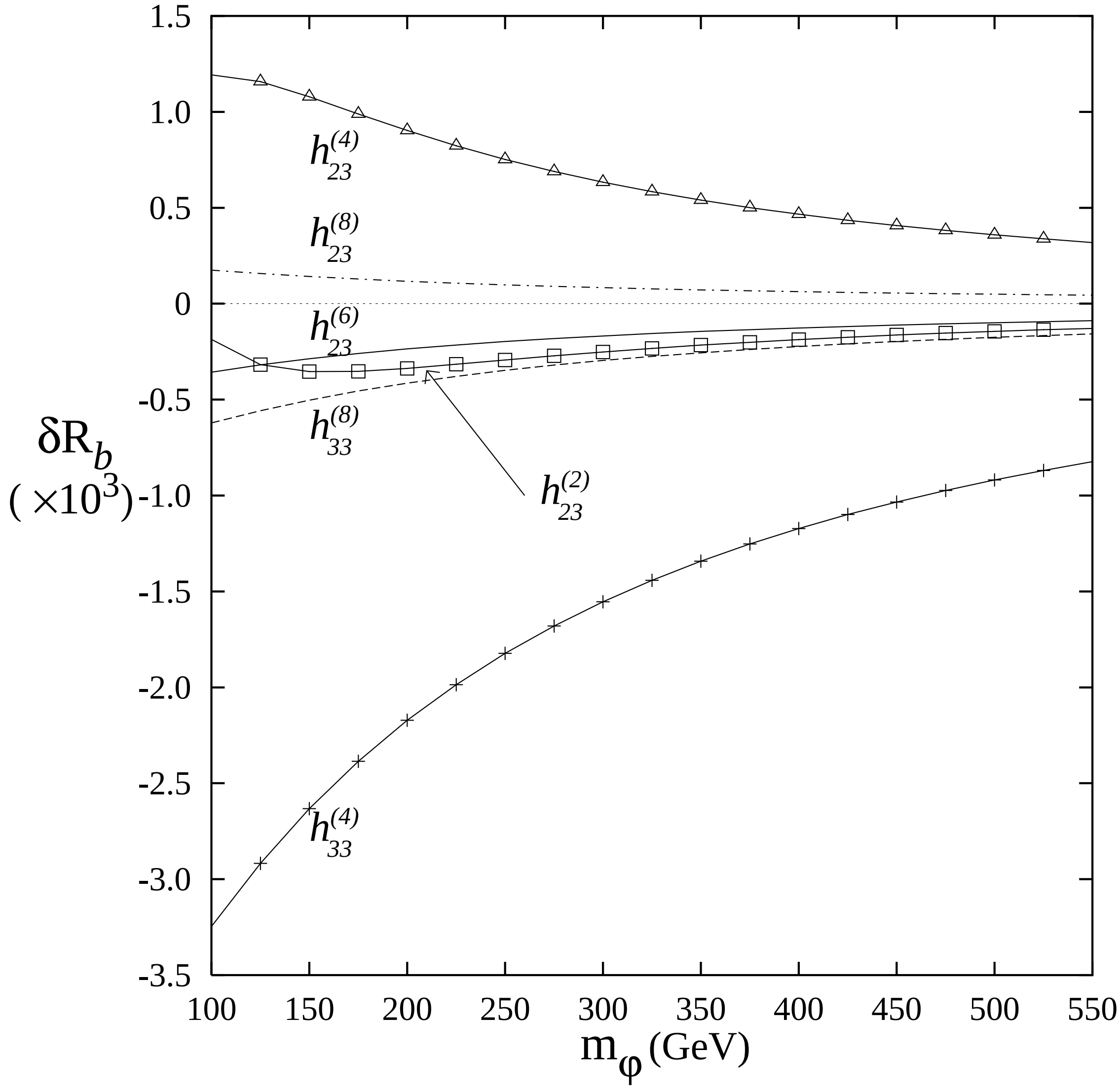

Figure 3: As in Fig. 2, but for R_b

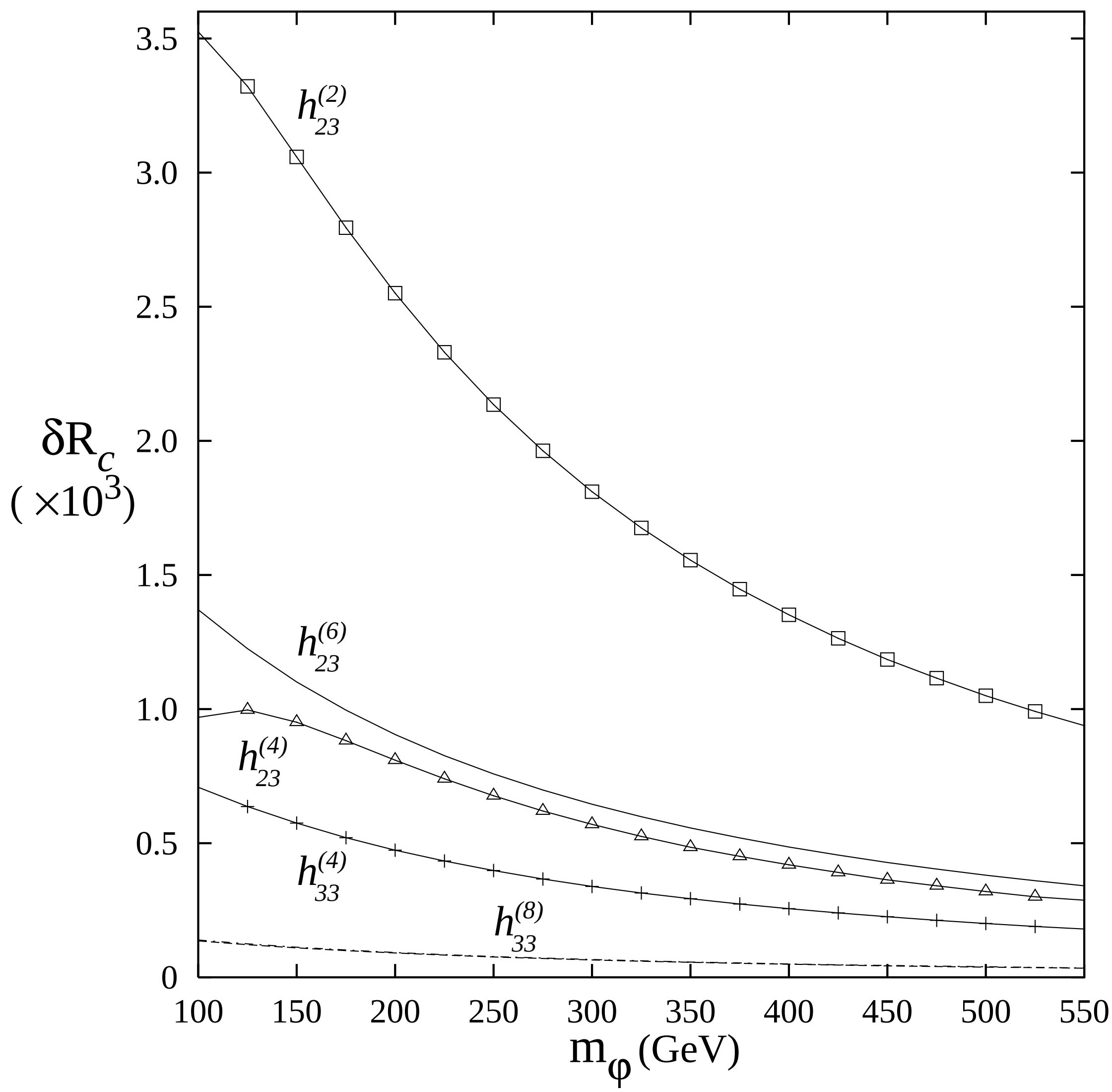

Figure 4: As in Fig.2, but for R_c