EUROPEAN ORGANIZATION FOR NUCLEAR RESEARCH

CERN-TH.5650/90

PSI-PR-90-06

Higher Order Dominance in the Standard Model

WEI-SHU HOU

*Paul Scherrer Institute
CH-5232 Villigen PSI, Switzerland*

and

ROBIN G. STUART

*Theory Division, CERN
CH-1211 Geneva 23, Switzerland*

ABSTRACT

Two examples are discussed where higher order processes dominate over lower order ones for dynamical reasons, as opposed to the usual condition of the existence of several expansion parameters. For QCD induced charmless B decays, $b \rightarrow sq\bar{q}$ dominates over $b \rightarrow sg$ due to the existence of two induced flavour changing gluonic couplings. In the b' system, the ' $\mathcal{O}(G_F^4)$ ', 3-body decay process $b' \rightarrow bZ^*$ dominates over the ' $\mathcal{O}(\alpha_s G_F^2)$ ', 2-body process because of different behavior under GIM cancellation, and because b' , if it exists, is not much lighter than M_Z . Confirmation or discovery of these phenomena will serve as tantalizing evidence for the inner workings of the Standard Model.

CERN-TH.5650/90

February, 1990

1. Introduction

In the usual application of perturbation theory, an expansion is made in some sufficiently small parameter assuming the convergence of the complete result. When there exist several small parameters in the model, even if perturbation theory can be consistently applied, there may exist special circumstances whereby, for a particular expansion parameter, higher order physical processes dominate. A good example is provided in the Standard Model (SM) by the so-called ‘altered’ Cabibbo patterns in B decays [1]. The decay mode $b \rightarrow s\bar{u}u$ is multiply-Cabibbo-suppressed in the SM, being proportional to $|V_{ub}V_{us}|^2$ as compared to $|V_{cb}|^2$ for the dominant $b \rightarrow c$ modes. QCD introduces new effective interactions at the one-loop order, namely the penguin diagrams, that suffer from usual loop suppression, amounting to a factor of $\alpha_s^2/(12\pi^2)$. However, they do not suffer from Cabibbo suppression with respect to the dominant $b \rightarrow c$ modes since the loop is dominated by c and t quarks and the Kobayashi-Maskawa (KM) matrix elements satisfy the condition, $V_{cs}^*V_{cb} \cong -V_{ts}^*V_{tb} \approx V_{cb}^*$. Because of a large logarithm in the effective coupling, and because of having three possible light quark species in the final state, the loop induced decay picks up a sizeable enhancement factor, and it turns out [2] that purely loop induced charmless $b \rightarrow s$ decays, $b \rightarrow s\bar{q}q$ (where $q = u, d, s$), have a branching ratio (BR) of order 1%. The current [3] crude experimental ‘measurement’ of $|V_{ub}/V_{cb}| \sim 0.1$ implies that the tree level contribution to $b \rightarrow s\bar{u}u$ is at the 0.1% level. Thus the conjecture, made ten years ago, of an ‘altered’ Cabibbo pattern is likely to be realized in Nature. We see that this is due to the combined effect of having two expansion parameters, V_{ub}/V_{cb} and α_s , for $b \rightarrow s$ decays without charmed particles in the final state.

An even more dramatic possibility was pointed out recently by us [4], in which higher order processes may even dominate over tree level processes and become the dominant decay modes. If a fourth generation exists, the b' may well be lighter than the top. Thus the decay $b' \rightarrow t$ is kinematically forbidden, while the $b' \rightarrow c$ term is probably suppressed by a small (compared to $V_{tb'}$) KM element $V_{cb'}$. A close analogy with the discussion in the previous paragraph can then be made. Note that current data indicate that charmless $b \rightarrow u$ (tree) and $b \rightarrow s$ (loop) decay rates are at similar order of magnitude. In a world where $m_b < m_c$, they would actually be the

dominant decay modes, since $b \rightarrow c$ would be kinematically forbidden. Furthermore, the tree level process is proportional to $|V_{ub}|^2$, hence in the case where $V_{cb} \rightarrow 0$, one would have a situation where loop-induced $b \rightarrow s$ transitions are the dominant B decay modes. We do not live in such a world, but we may live in an analogous world if the fourth generation exists. If $m_{b'} \lesssim m_t$, and if $V_{cb'}$ is sufficiently small compared to $V_{tb'}$, then loop-induced $b' \rightarrow b$ decays may dominate the total b' rate rather than just be dominant over some specific lower order process. The consequences are not just theoretical. If the b' rate is dominated by flavour-changing neutral current processes, the usual charged current signature of “isolated $e/\mu + \text{jets}$ ” for heavy quark production may be absent, and different search strategies have to be devised [4].

Intriguing as they are, the previous examples constitute the ‘trivial’ kind of higher order dominance. It is just due to a combination of kinematical conditions and the existence of multiple expansion parameters. Of course, the loop amplitudes themselves require some dynamics that does not suppress them completely. In this paper we wish to discuss further two cases that are even more remarkable, namely $b \rightarrow s\bar{q}q \gg b \rightarrow sg$ and $b' \rightarrow bZ^* \gtrsim b' \rightarrow bg$ [5], in which the higher order process dominates for purely dynamical reasons. Both cases involve processes that have the same KM structure. In the first, B physics case, the $\mathcal{O}(\alpha_s^2)$ process is found to dominate over the $\mathcal{O}(\alpha_s)$ process. In the second case, the $\mathcal{O}(G_F^4)$ process is found to dominate over the $\mathcal{O}(\alpha_s G_F^2)$ process. In both cases, the dominant process even suffers further from phase space suppression. It is our purpose here to elucidate the underlying physics that leads to such peculiar results.

2. The Higher Order Dominance, $b \rightarrow s\bar{q}q \gg b \rightarrow sg$

Let us start with QCD induced charmless B decays. The effective $s\bar{b}g$ coupling induced at one-loop order via charged current interactions is of the form demanded by current conservation,

$$\Gamma_\mu^a = g_s \frac{G_F}{4\sqrt{2}\pi^2} V_{ts}^* V_{tb} \bar{s} t^a \{ F_1^i(q^2 \gamma_\mu - q_\mu \not{q}) L - F_2^i \varepsilon_{\mu\nu\alpha\beta} q^\nu m_b R \} b, \quad (1)$$

where g_s is the strong coupling constant, $t^a \equiv \lambda^a/2$ and λ^a is the Gell-Mann matrix, $q = p_b - p_s$. Note that there are two effective couplings and in what follows we will

refer to them as form factors although q^2 dependences will be ignored for the present discussion. The ‘charge radius’ form factor for each intermediate quark $i = u, c, t$, after subtraction of an arbitrary constant permitted by the GIM mechanism [6], is

$$F_1^i = \frac{x_i}{12} \left(\frac{1}{y_i} + \frac{13}{2} \frac{6}{y_i^2} - \frac{6}{y_i^3} \right) + \left[\frac{2}{3y_i} - \frac{x_i}{6} \left(\frac{4}{y_i^2} + \frac{5}{y_i^3} - \frac{3}{y_i^4} \right) \right] \log x_i, \quad (2)$$

while the ‘dipole moment’ form factor is

$$F_2^i = -\frac{x_i}{4} \left(-\frac{1}{y_i} + \frac{3}{y_i^2} + \frac{6}{y_i^3} \right) + \frac{3x_i}{2y_i^4} \log x_i, \quad (3)$$

where $x_i = m_i^2/M_W^2$ and $y_i = x_i - 1$. These results were calculated by expanding to leading non-trivial order in external masses and momenta while keeping the exact dependence on internal masses.

The important point is the presence of two effective couplings [2]. If only one effective coupling existed, then one would normally have the usual perturbative result of $b \rightarrow sg \gg b \rightarrow sg^*$, where in the latter, the gluon is off-shell (see Fig. 1). This would just be the naïve parton model result, since it is both lower order in α_s , and phase space favoured. However, with two effective couplings, one notices that for the on-shell gluon $b \rightarrow sg$ decay, only a dipole moment transition is possible, since the charge radius transition vanishes with q^2 . Thus, only the F_2 effective coupling contributes to the $b \rightarrow sg$ transition which, as argued, is presumably the leading one. If for some reason the F_1 effective coupling is much stronger than the F_2 coupling, the possibility emerges that $b \rightarrow sg^*$ would dominate over the lower order ‘parton model’ process $b \rightarrow sg$. This is precisely what happens [2] in the B system, given the known parameters within SM.

Since^[1] $V_{ub}V_{us}$ is rather small compared to $V_{cb}V_{cs} \cong -V_{tb}V_{ts}$, loop induced $b \rightarrow s$ transitions are dominated by internal c and t quark effects. Notice from eqs. (2) and (3), that for small x_i ,

$$F_2^i \cong \frac{1}{2} x_i, \quad (4a)$$

whereas,

$$F_1^i \cong -\frac{2}{3} \log x_i, \quad (4b)$$

[1] We ignore the possibility of CP violation and assume real KM elements in the following discussion.

while both functions become rather damped and tend to a constant ($F_1^i \rightarrow 1/12$, $F_2^i \rightarrow 1/4$) for large x_i . Eq. (4b) indicates infrared sensitivity as the internal quark mass becomes small, which is naturally regulated by q^2/M_W^2 in the infrared, as a more complete treatment [7] would show. It is clear, however, that ‘light’ (as compared to m_b) internal quark effects are sizeable for F_1 , although quite negligible for F_2 . The physical interpretation is the following. Treating m_c as a variable mass, when it becomes sufficiently light, there is a large direct branching ratio into $b \rightarrow c\bar{c}s$ final state. The colour octet, ‘on-shell’ $c\bar{c}$ configuration has a substantial rescattering probability into charmless $q\bar{q}$ final states which is reflected in the large logarithm for the charge radius form factor (eq. (4b)). However, for the dipole transition, a spin flip is required, and given the absence of large fermion masses for a light fermion final state, the rescattering amplitude is damped by powers of m_c^2/M_W^2 (eq. (4a)). From a different perspective, one can say that GIM cancellation leads only to a logarithmic suppression in internal quark masses for the charge radius form factor F_1 , because of the necessity of a q^2 factor from gauge invariance, while for F_2 , the spin flip condition requires a power suppression in internal quark masses. The upshot is that, with m_c taken as 1.5 GeV, F_1^c is roughly 5 while F_2^c is of order 10^{-4} . For $m_t \sim 80$ GeV, $F_1^t \approx 0.6$ and F_2^t is five times smaller, but as m_t becomes very large, $F_1^t \rightarrow 1/12$ whereas $F_2^t \rightarrow 1/4$.

Taking $m_c = 1.5$ GeV, $m_b \approx 5$ GeV, $\alpha_s(m_b^2) \cong 0.23$, we plot the SM result for charmless $b \rightarrow s$ decays in Fig. 2. The space-like $b\bar{q} \rightarrow s\bar{q}$ transition is incorporated into $b \rightarrow sg^*$, and anyway constitutes only a small fraction because of the smallness of $(f_B/m_B)^2$. For $b \rightarrow sg$, both the ‘bare’ (purely electroweak loop) and QCD corrected results [8] are plotted. Note that $b \rightarrow sg^*$ and the ‘bare’ $b \rightarrow sg$ result rise with m_t for different reasons. The decay, $b \rightarrow sg^*$ is dominated by ‘long distance’ (compared to $1/m_b$) charm contributions, and the t quark contribution is destructive due to GIM cancellation. The decay rate rises with m_t because F_1^t drops with m_t . For the ‘bare’ $b \rightarrow sg$, however, the t contribution is leading, which grows with m_t . QCD corrections tend to further suppress $b \rightarrow sg$.

This last point is in stark contrast to another well-known case of ‘higher order dominance’. In the analogous case of $b \rightarrow s\gamma$, the photonic F_2 form factor suffers from a similar power GIM cancellation, which leads to a rather suppressed BR for light m_t . However, unlike the $b \rightarrow sg$ case, QCD corrections actually dominate

3. The Higher Order Dominance, $b' \rightarrow bZ^* \gtrsim b' \rightarrow bg$

The situation is different as one turns to the fourth generation^[2] down-type quark, b' . As mentioned earlier, the b' system affords us the dramatic possibility of having loop induced $b' \rightarrow b$ transitions dominate the total decay rate, thereby leading to the possibility of discovering the b' quark through its FCNC decay process [4]. The arguments have been given earlier, and the result after extensive computation, is that, for $m_{b'} \lesssim M_W$, $M_W \lesssim m_{b'} \lesssim M_Z$ and $m_{b'} \gtrsim M_Z$, the decay $b' \rightarrow b$ may dominate over $b' \rightarrow c$ if $V_{cb'}/V_{ub'} \lesssim 10^{-2}$, 10^{-3} and 10^{-2} , respectively. For $m_{b'}$ between M_W and M_Z , the charged current $b' \rightarrow cW$ decay leads to an on-shell W , with corresponding 2-body phase space, and therefore a smaller $V_{cb'}$ is needed to suppress it below $b' \rightarrow b$. However once $m_{b'}$ reaches a value close to M_Z or higher, the decay $b' \rightarrow bZ$ turns on, and the situation is restored. Below the M_Z threshold, $b' \rightarrow bg$ and $b' \rightarrow bZ^*$ ^[3] are on an equal footing, with $b' \rightarrow bZ^*$ tending to dominate. The useful signature of $b' \rightarrow b\gamma$ is at the 10% level.

Before we discuss the bizarre situation of $b' \rightarrow bZ^* \gtrsim b' \rightarrow bg$, it is interesting to compare $b' \rightarrow bg$ vs. $b' \rightarrow bZ^*$ to the B meson case. The b' is a fourth generation quark and given the likely hierarchy of $V_{cb'} \ll V_{ub'}$, its loop induced decay will be dominated by internal t and t' quarks. Thus, only the third and fourth generation quarks participate in $b' \rightarrow b$ transitions. Unlike the b quark, $m_{b'}$ is of order M_W , and in fact, $m_{b'} \gtrsim M_W$ as well, so no large logarithms appear in the loop integral. Although there are now five light quark species accessible, $\alpha_s(m_{b'})$ is smaller. Given that there are no more large logarithms, it turns out that F_1 and F_2 are now comparable, and one has the usual perturbative result of $b' \rightarrow bg \gg b' \rightarrow bZ^*$ [4]. Thus the naive parton model works, but still for non-trivial reasons. This also highlights the peculiarity of the B system discussed above.

The comparison between $b' \rightarrow bg$ and $b' \rightarrow bZ^*$ is all the more surprising. In

^[2]We do not wish to justify the existence of a fourth family, but rather, to concentrate on the peculiar features that it may offer. Present data only excludes the existence of a fourth light neutrino species.

^[3]In this note we use only the virtual Z contribution for sake of simplicity. Gauge invariance demands the inclusion of box diagrams and virtual photon contributions, but the qualitative features are not changed. For discussions see Ref. 5.

over the 'bare' contribution for all practical top mass values [9], for $m_t \lesssim 100$ GeV, substantially so. The difference with the $b \rightarrow sg$ case can be understood from the lowest order result. QCD corrections bring in large logarithms as $\mathcal{O}(\alpha_s)$ corrections to the 'bare' $\bar{s}b\gamma$ and $\bar{s}bg$ couplings. For the former, the sign is positive, and the coefficient is large, while for the latter the sign is opposite, with a rather small coefficient. Thus, $b \rightarrow s\gamma$ is greatly enhanced by QCD corrections, whereas for $b \rightarrow sg$, QCD corrections not only do not dominate the rate, it actually suppresses it by cancellation. Thus, the 'large' QCD enhancement of $b \rightarrow s\gamma$ is another good example of higher order dominance. However, note that it is not the QCD correction itself that is large, but rather the dynamical power GIM cancellation of the dipole form factor that suppresses the leading order electroweak loop result. In this way, $b \rightarrow s\gamma$ is on the same footing as our example of $b \rightarrow sg^*$ dominating over $b \rightarrow sg$, although the latter is more intricate. For $b \rightarrow sg^*$, since a large logarithm already appears at leading order, QCD corrections do not have a pronounced effect.

The dominance of $b \rightarrow sg^*$ over $b \rightarrow sg$ is unique to the B system, due to a combination of several factors. First of all, the b quark is of the third generation, and for its loop induced decay c and t contributions are equally important. Although of the third generation, its mass of 5 GeV is such that it is very small compared to M_W , while large compared to m_c , viz. $m_c^2 \ll m_b^2 \ll M_W^2$, and thus allows large logarithms to appear from the loop integral. Finally, $\alpha_s(m_b^2)$ is not small, therefore there is not a huge price to pay in going to the higher order $b \rightarrow sg^*$ process. Having three light quarks accessible in the final state compensates for the extra power of α_s , and one has only to compensate for the loss in phase space. As we have seen, the subtle difference in GIM cancellation and the fact that only the highly GIM suppressed dipole form factor contributes to the lower order process, combined with all the above-mentioned features, allowed the higher order process $b \rightarrow sg^*$ to dominate.

Fig. 3 we give typical diagrams that give rise to these processes. Note that, $b' \rightarrow bg$ is of naïve order $\mathcal{O}(\alpha_s G_F^2)$ with 2-body phase space, while $b' \rightarrow bZ^*$ is of naïve order $\mathcal{O}(G_F^4)$, with 3-body phase space. On both counts, especially G_F^2 vs. G_F^4 , one would conclude that $b' \rightarrow bg$ would by far dominate. In fact, it has been stated in the literature [10] that “the off-shell Z^0 contribution is one order of magnitude or more below the $b' \rightarrow b\gamma$ transition”, which would put it two orders of magnitude below $b' \rightarrow bg$, and illustrates how very peculiar it sounds to have $b' \rightarrow bZ^* \gtrsim b' \rightarrow bg$. Direct computation [5] bears out that this is indeed the case, as shown in Fig. 4. It is therefore of some interest to understand the cause.

The effective $\bar{b}b'Z$ coupling induced at the electroweak one-loop level is of the form

$$\Gamma_\mu = \frac{g}{\cos\theta_W} \frac{G_F}{2\sqrt{2}\pi^2} V_{ib}^* V_{ib'} \bar{b}(p) [\gamma_\mu F_L^i L - i\sigma_{\mu\nu} \frac{q_\nu}{M_Z} F_{TR}^i R] b'(p'), \quad (5)$$

where $q = p_W - p_b$, and we have factored out relevant factors so one can directly compare with the effective $\bar{b}b'g$ and $\bar{b}b'\gamma$ vertices, which are of the form shown in eq. (1). Note that in the case of the massive vector boson Z^0 , a natural scale, M_Z , appears which is used to scale the explicit power of q for the dipole term. Also, since massive vector gauge bosons are not associated with conserved currents, the ‘charge radius’ form factor can now be more general, not necessarily vanishing as $q^2 \rightarrow 0$. A massive boson can couple to a non-vanishing ‘charge radius’ at long distances, since the short range nature of the boson propagator will damp it out anyway. This turns out to be the crucial element.

The important point to note is that^[4] F_L has dimension M^2 as can be seen by comparing eq. (5) with eq. (1). The $\sin^2\theta_W$ dependent part of the effective coupling is, of course, conserved and a q^2 factor can be extracted just as in the gluonic and photonic case. However, the remainder of the effective Z coupling couples to a non-conserved current, and an explicit q^2 factor is not necessary. We have argued earlier that the need for an explicit q^2 factor allows for a logarithmic type of GIM cancellation. As seen from eqs. (2) and (3), for large internal quark masses, the effective F_1 and F_2 couplings to the gluon (or photon) are damped by inverse

^[4] F_{TR} turns out to be subleading compared to F_L , and we therefore leave it out of our discussion.

powers of $(m_t^2/M_W^2 - 1)$. Thus, current conservation, which is intimately related to unbroken-gauge symmetries and massless gauge bosons, seems to ‘screen out’ heavy internal fermion mass effects, thereby more or less satisfying the decoupling theorem [11], albeit in a rather subtle way. The decoupling theorem is evaded in the broken sector, when currents are not conserved, and their divergence related to densities via explicit powers of masses. This is not all that surprising, since the fermions derive their masses via Yukawa couplings to the Higgs field which is a dynamical effect. The Goldstone bosons generated by spontaneous symmetry breaking get eaten and become the ‘longitudinal component’ of the W and Z fields, and the symmetry breaking scale (vacuum expectation value) is responsible for both the intermediate vector boson masses and the fermion masses. In this way, the W and Z bosons ‘know’ more about the fact that the symmetry is broken, and that the fermion masses are generated dynamically in relation to symmetry breaking.

What we wish to emphasize is the following. The leading dependence of the form factor F_L^i is proportional to the internal fermion mass squared, m_t^2 , regardless of whether the external masses and momenta are kept. This is not particularly surprising but rather it is a non-trivial manifestation of the symmetry breaking and mass generation mechanism (for both bosons and fermions) of the Standard Model.

The phenomenological consequences are intriguing. We had said that $b' \rightarrow bZ^*$ is naively of order $\mathcal{O}(G_F^4)$, whereas $b' \rightarrow bg$ is of order $\mathcal{O}(\alpha_s G_F^2)$. One could have noticed that the dimensionality differed by M^{-4} , and the foregoing arguments suggest that it is compensated roughly by $(m_t^2 - m_b^2)^2$. Given that m_t is presumably rather heavy, this factor effectively removes a factor of G_F^2 . What is more, the rather weak asymptotic dependence of F_L^i on internal quark masses (weaker than logarithmic, and tending towards a constant), together with the fact that m_t is taken to be rather heavy already, lead to severe GIM cancellation for the effective $\bar{b}b'g$ coupling. For $\bar{b}b'Z$ the GIM cancellation is minimal so long as m_t is not too close to m_b . This forms the basis of the peculiar case of higher order dominance, $b' \rightarrow bZ^* \gtrsim b' \rightarrow bg$. Even the unfavourable 3-body phase space is overcome.

There is one final point that helps considerably towards overcoming phase space suppression. Normally, combining phase space and matrix-element-squared there is a suppression factor of $(m_W^2 - q^2)^2$ in the 3-body differential rate. However, in the

case of the decay $b' \rightarrow bZ'$, there is a propagator factor $(q^2 - M_Z^2)^{-2}$. Given that $m_{b'}$ is not that small compared to M_Z , unlike the B meson case, these two factors almost compensate each other and the differential rate is rather flat in q^2 up to rather high values before the first suppression factor sets in [5]. The situation is not different from say $t \rightarrow bW^+$ decay, and is usually called the “massive gauge boson propagator effect” [12]. It compensates rather well for the 3-body phase space. In comparison, in the B meson case, since $m_b^2 \ll M_W^2$, the b quark essentially goes through a four-fermion interaction induced decay. This, together with the peculiar combination that enhances $b \rightarrow sg^*$, makes $b \rightarrow sZ'$ in comparison 2–3 orders of magnitude smaller within the Standard Model; what we said earlier regarding m_t^2 enhancement of the $\bar{s}bZ$ vertex still holds true though [13]. However, $b \rightarrow sZ'$ does dominate over $b \rightarrow s\gamma'$ so long as m_t is large enough, which is again some form of higher order dominance.

4. Conclusion

We have illustrated two particular examples of higher order dominance, which are not just due to the existence of several perturbative parameters, but rather, dynamical effects in the loop amplitude. In both cases, the effect is related to how the GIM cancellation gets implemented in the loop process. There are basically three kinds of GIM cancellation in the Standard Model: power and logarithmic GIM that are damped by inverse powers of internal fermion mass-squared when they become heavy, and power GIM cancellation that manifestly evades the decoupling theorem. Although it is the symmetry breaking and mass generation mechanism itself that permitted incomplete GIM cancellation, whether the specific external gauge boson couples to conserved currents or not makes a rather big difference. Massless gauge bosons, corresponding to unbroken symmetries, have loop induced effective couplings that almost satisfy the decoupling theorem, i.e. the first two types of GIM cancellation. Although generating some effects for internal fermion masses that are below roughly M_W , for heavier internal fermions, the effects are rather ‘decoupled’. In this respect, the induced gluonic FCNC coupling is even more decoupled from internal heavy fermion effects than the photonic one, since QCD is ‘unaware’ of electroweak symmetry breaking, while the photon is ‘aware’.

as evidenced by the W - W - γ coupling and gives a slightly stronger dependence on heavy internal fermion masses. For the massive Z boson, which couples to a non-conserved current corresponding to a broken symmetry, its loop-induced effective couplings satisfy the third kind of GIM cancellation, i.e. those of the power type, which are strongly dependent on and rise rapidly with heavy internal fermion masses. They are rather effective probes of virtual, heavy particle effects, since the Z boson itself is intimately involved in the highly intertwined mechanism of gauge symmetry breaking, gauge boson mass generation, and fermion mass generation, within the Standard Model framework. Of course, the effective FCNC Higgs coupling shares this same feature.

In summary, $b' \rightarrow bg^*$ dominates over $b' \rightarrow bg$ because logarithmic GIM cancellation is operative for the charge radius form factor, F_1 , and power GIM for the dipole moment form factor, F_2 . Only the latter contributes to the on-shell $b \rightarrow sg$ transition. A unique situation in the B system leads to a large logarithm in F_1 which results in this unusual case of higher order dominance. Even more intriguing is the case $b' \rightarrow bZ'^+ \gtrsim b' \rightarrow bg$, where the $\mathcal{O}(G_F^4)$, 3-body process dominates over the $\mathcal{O}(\alpha_s G_F^2)$, 2-body process. This is especially so if $m_{t'}$ is very heavy. The non-decoupled $m_{t'}$ ² dependence of the effective F_L coupling for $b' \rightarrow bZ'$, together with the absence of large logarithms and the almost decoupled nature of the F_2 coupling of $b' \rightarrow bq$, give rise to this very peculiar result. The examples of higher order dominance presented here follow from the intricate nature of symmetry breaking and mass generation mechanism for both fermions and bosons in the Standard Model. Thus, the observation or discovery of such phenomena would be of great interest in shedding light on the underlying mechanism.

FIGURE CAPTIONS

1. Typical diagrams contributing to (a) $b \rightarrow sg$ and (b) $b \rightarrow sg'$.
2. Branching ratio for QCD induced $b \rightarrow s$ transitions. Solid for $b \rightarrow sg'$, dashed for $b \rightarrow sg$ without QCD corrections, dot-dashed for QCD corrected $b \rightarrow sg$.
3. Typical diagrams contributing to (a) $b' \rightarrow bg$ and (b) $b' \rightarrow bZ^*$.
4. Decay rate for leading $b' \rightarrow b$ transitions: (a) $b' \rightarrow bg$, (b) $b' \rightarrow bZ^*$. The KM factor $|V_b^* V_{t'b'}|^2 \approx |V_{tb}^* V_{t'b'}|^2$ has been removed. The dashed curves are for $m_\mu = 300$ GeV and the solid curves are for $m_\mu = 200$ GeV, while $m_\mu = 80, 120$ GeV for upper and lower curves, respectively, in each set.

REFERENCES

- [1] R. Guberina, R.D. Peccei and R. Rückl, *Phys. Lett.* **90 B** (1980) 169.
- [2] W.S. Hou, *Nucl. Phys.* **B 308** (1988) 561.
- [3] R. Fulton *et al.* (CLEO Collaboration), *Phys. Rev. Lett.* **64** (1990) 16; H. Albrecht *et al.* (ARGUS Collaboration), *Phys. Lett.* **B 234** (1990) 409.
- [4] W. S. Hou and R. G. Stuart, *Phys. Rev. Lett.* **62** (1989) 617; *Nucl. Phys.* **B 320** (1989) 277.
- [5] W. S. Hou and R. G. Stuart, preprint CERN TH.5601/89.
- [6] S. Glashow, J. Iliopoulos and L. Maiani, *Phys. Rev. D* **2** (1970) 1285
- [7] See, e.g. J.-M. Gérard and W.S. Hou, *Phys. Rev. Lett.* **62** (1989) 855; W.S. Hou and R.S. Willey, *Nucl. Phys.* **B 326** (1989) 54.
- [8] R. Grigjanis, P. J. O'Donnell, M. Sutherland and H. Navalet, *Phys. Lett.* **B 224** (1989) 209.
- [9] S. Bertolini, F. Borzumali and A. Masiero, *Phys. Rev. Lett.* **59** (1987) 180; N. G. Deshpande, P. Lo, J. Trampetic, G. Ellam and P. Singer, *Phys. Rev. Lett.* **59** (1987) 184; B. Grinstein, R. Springer and M. Wise, *Phys. Lett.* **B 202** (1988) 138; R. Grigjanis P. J. O'Donnell, M. Sutherland and H. Navalet, *Phys. Lett.* **B 213** (1989) 355.
- [10] K. Abe *et al.* (VENUS Collaboration), *Phys. Lett.* **B 234** (1990) 382.
- [11] T. Appelquist and J. Carazzone, *Phys. Rev. D* **11** (1975) 2856.
- [12] I. Bigi, Yu. I. Dokshitzer, V. Khoze, J. H. Kühn, and P. Zerwas, *Phys. Lett.* **B 181** (1986) 157.
- [13] W. S. Hou, R. S. Willey and A. Soni, *Phys. Rev. Lett.* **58** (1987) 1608.

Fig. 1

Fig. 2

Fig. 3

Fig. 4