A Possible Arena for Searching New Physics - the $\Gamma(D^0 \to \rho^0\gamma)/\Gamma(D^0 \to \omega\gamma)$ Ratio

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ABSTRACT

We propose to investigate flavour changing neutral currents in the $c \to u\gamma$ transition through the measurement of the difference between $\Gamma(D^0 \to \rho^0\gamma)$ and $\Gamma(D^0 \to \omega\gamma)$. This is based on the observation that $D^0 \to (d\bar{d})\gamma$ is due to long distance physics while $D^0 \to (u\bar{u})\gamma$ arises from the $c \to u\gamma$ transition. The effect of $\rho - \omega$ mixing is included. A difference in the decay widths of more than about 30% would be indicative of new physics.

Despite the remarkable success of the standard model (SM), it is generally believed that this is an effective theory at present energies. The lack of explanation for many of the salient features of SM suggests that one must look for its extension. At present, there is no clear picture of the form of the SM extension and the search for physics beyond the SM proceeds along many alternative paths.

As the couplings of Z^0, photon and Higgs boson are flavour diagonal, flavour changing neutral current (FCNC) transitions are rare in SM since they can arise at loop level only. Moreover, FCNC transitions in the up-quark sector (like $c \to u\gamma$) are much rarer than those in the down-quark sector (like $b \to s\gamma$) as a result of the CKM matrix elements and masses involved [1]. Accordingly, they could provide an appropriate ground for the search of new physics effects [2]. In fact, rare charm decays have been frequently marked as possible sources for the discovery of new physics [2, 3], in view of the smallness of the short distance SM contributions of these processes.

Hadronic decays like $D \to V\gamma$ [1, 4-9] and $D \to Vl^+l^-$ [9-12] have been considered with the aim of investigating the size of the short distance (SD) $c \to u\gamma$, $c \to ul^+l^-$ contribution to the decay amplitude. However, it turns out that these decays are largely dominated by long distance (LD) contributions, rendering them inappropriate for detecting deviations from the SM values of the basic $c \to u\gamma(l^+l^-)$ transitions. An exception is $B_c \to B_u^*\gamma$ since in this decay both the SD and LD contributions to the branching ratio are in the 10^{-8} range [13].

In the present letter we suggest a new possibility for the search of the $c \to u\gamma$ transition, from the determination of the difference in size of the partial decay widths $D^0 \to \rho^0\gamma$ and $D^0 \to \omega\gamma$. Our method is particularly suitable for detecting an enhancement of the $c \to u\gamma$ amplitude coming from new physics if it increases this amplitude by at least a factor 3-4. We remark that enhancements of up to a factor 100 have been noted in certain non-minimal[1] realizations of supersymmetry [2].

Next we note that in the radiative decays of D^0 mesons, the $(d\bar{d})\gamma$ final state arises primarily as a result of nonleptonic W- exchange $c\bar{u} \to d\bar{d}$, being therefore an outcome of LD physics (Fig 1a), while the $(u\bar{u})\gamma$ final state is mainly due to the electro-magnetic penguin $c \to u\gamma$ transition (Fig. 1b). The $(u\bar{u})\gamma$ final state receives also a small contribution form the long distance penguin mechanism illustrated in Fig. 1c. The decays $D^0 \to (\rho, \omega)\gamma$ have been treated in great detail in Refs. [7, 8, 9] and it is shown that the LD contributions completely overshadow the SD contribution in SM. Although, there is a strong enhancement at the two - loop level [14] of the $c \to u\gamma$ amplitude [15] as a result of gluonic corrections, the QCD - corrected SM amplitude [14] is still only a few percent of the LD one [8].

As usual, we define the isospin eigenstates

$$\omega^{(I=0)} = \frac{1}{\sqrt{2}}(\bar{u}u + \bar{d}d) \ , \qquad \rho^{(I=1)} = \frac{1}{\sqrt{2}}(\bar{u}u - \bar{d}d) \ . \tag{1}$$

Therefore, in the absence of the $\omega - \rho$ mixing and the LD penguin contribution, the LD rates for $D^0 \to \rho\gamma$ and $D^0 \to \omega\gamma$ are equal in the SM; a difference must be due to SD effects. This is the basis for our approach.

The main decay modes of the physical ρ, ω states are 2π and 3π respectively. It is well known that the $\omega^{(I=0)}$ and $\rho^{(I=1)}$ states can mix and ω is known to decay also to 2π states, with a branching ratio of $Br(\omega \to \pi^+\pi^-) = (2.21 \pm 0.30)\%$ [16]. The physical states denoted as ρ and ω are related to the states defined in Eq. (1) by [17]

$$\rho = \rho^{(I=1)} - \epsilon\ \omega^{(I=0)} \ , \qquad \omega = \omega^{(I=0)} + \epsilon\ \rho^{(I=1)} \ , \tag{2}$$

where

$$\epsilon = \frac{\Pi^2_{\rho\omega}}{\hat{m}^2_\omega - \hat{m}^2_\rho} \simeq \frac{\Pi^2_{\rho\omega}}{m^2_\omega - m^2_\rho + im_\rho\Gamma_\rho - im_\omega\Gamma_\omega} \tag{3}$$

with complex masses $\hat{m} = m + im\Gamma/2$ and with real m and Γ. This leads to the $\omega \to 2\pi$ amplitude

$$\mathcal{A}(\omega \to 2\pi) = \frac{\Pi^2_{\rho\omega}}{m^2_\omega - m^2_\rho + im_\rho\Gamma_\rho - im_\omega\Gamma_\omega}\ \mathcal{A}(\rho \to 2\pi), \tag{4}$$

from which the mixing parameter is determined to be

$$\Pi^2_{\rho\omega} \ = \ -(4.0 \pm 0.4) \times 10^{-3}\ \ \text{GeV}^2 \tag{5}$$

[1]Non-minimality here means that the universality of the soft breaking terms is not imposed or that additional super-fields are added to MSSM.

where the values of m_ρ, m_ω and Γ_ρ are taken from [16]. The minus sign in (5) is obtained from a detailed analysis of the $e^+e^- \to \pi^+\pi^-$ amplitude [17]. The corresponding value of ϵ (3) is

$$\epsilon \;=\; -0.0061 + i\;0.036, \tag{6}$$

and the physical states can be written as

$$\rho = \frac{1}{\sqrt{2}}(1-\epsilon)u\bar{u} + \frac{1}{\sqrt{2}}(-1-\epsilon)d\bar{d} \qquad \omega = \frac{1}{\sqrt{2}}(1+\epsilon)u\bar{u} + \frac{1}{\sqrt{2}}(1-\epsilon)d\bar{d}\;. \tag{7}$$

The general form of a $D \to V\gamma$ transition amplitude is written as usual in terms of parity conserving A_{PC} and parity violating A_{PV} amplitudes [1, 7, 8]

$$\begin{aligned}\mathcal{A}[D^0(p) \to V(p_V,\epsilon_V)\gamma(q,\epsilon_\gamma)] = -i\{&A_{PC}\epsilon_{\mu\nu\alpha\beta}q^\mu\epsilon_\gamma^{*\nu}p^\alpha\epsilon_V^{*\beta} \\ &+iA_{PV}[(\epsilon_V^*\cdot q)(\epsilon_\gamma^*\cdot p_V)-(q\cdot p_V)(\epsilon_V^*\cdot\epsilon_\gamma^*)]\}.\end{aligned} \tag{8}$$

We decompose the A_{PC} and A_{PV} amplitudes according to their final state content

$$A_{PC} = A_{PC}(u\bar{u}) + A_{PC}(d\bar{d}), \qquad A_{PV} = A_{PV}(u\bar{u}) + A_{PV}(d\bar{d})\;. \tag{9}$$

At this point we define the new quantities η_{PC}, η_{PV}

$$\eta_{PC} = \frac{A_{PC}(u\bar{u})}{A_{PC}(d\bar{d})}\;, \qquad \eta_{PV} = \frac{A_{PV}(u\bar{u})}{A_{PV}(d\bar{d})}\;, \tag{10}$$

which, as we explain later on, are of the order of a few percent in the standard model, and the amplitudes are rewritten as

$$\begin{aligned}A_{PC/PV}(D^0 \to \rho^0\gamma) &= \frac{1}{\sqrt{2}}A_{PC/PV}(d\bar{d})\;[(1-\epsilon)\eta_{PC/PV} + (-1-\epsilon)]\;,\\ A_{PC/PV}(D^0 \to \omega\gamma) &= \frac{1}{\sqrt{2}}A_{PC/PV}(d\bar{d})\;[(1+\epsilon)\eta_{PC/PV} + (1-\epsilon)].\end{aligned} \tag{11}$$

The decay width is given by

$$\Gamma(D \to V\gamma) = \frac{1}{4\pi}(\frac{m_D^2 - m_V^2}{2m_D})^3(|A_{PC}|^2 + |A_{PV}|^2)\;. \tag{12}$$

In order to extract the $(u\bar{u})\gamma$ final state, we propose a quantity $\mathcal{D}^{\omega-\rho}$, defined as

$$\mathcal{D}^{\omega-\rho} \equiv \frac{\Gamma[D^0 \to \omega\gamma]/(m_D^2 - m_\omega^2)^3 - \Gamma[D^0 \to \rho^0\gamma]/(m_D^2 - m_\rho^2)^3}{\Gamma[D^0 \to \omega\gamma]/(m_D^2 - m_\omega^2)^3}\;. \tag{13}$$

The standard model values of $\eta_{PC/PV}$ (given in (20) bellow) and ϵ (6) are small and $\mathcal{D}^{\omega-\rho}$ can be expanded to the first order

$$\mathcal{D}^{\omega-\rho} \simeq 4\;\frac{|A_{PC}(d\bar{d})|^2\;(\mathrm{Re}\;\eta_{PC} - \mathrm{Re}\;\epsilon) + |A_{PV}(d\bar{d})|^2\;(\mathrm{Re}\;\eta_{PV} - \mathrm{Re}\;\epsilon)}{|A_{PC}(d\bar{d})|^2 + |A_{PV}(d\bar{d})|^2}\;, \tag{14}$$

which leads to the near equality of the $D^0 \to \rho\gamma$ and $D^0 \to \omega\gamma$ rates. On the other hand, physics beyond the standard model, which affects significantly the size of $c \to u\gamma$ transition, will lead to a splitting of the $\Gamma[D^0 \to \rho^0\gamma] \simeq \Gamma[D^0 \to \omega\gamma]$ near equality. Sizeable coefficients $\mathcal{D}^{\omega-\rho}$ can arise in such scenarios and the expansion (14) is not valid. We remark that as a consequence of the mixing term $\epsilon \neq 0$ (2) there is a difference of a few percent between the widths of the two decays, even when η_{PC} and η_{PV} are neglected. An additional difference arises at the hadronic level as a consequence of $SU(3)$ flavour breaking. For example, as one sees from the explicit expressions of [7], the difference in couplings and masses of ρ and ω induces a $\sim 10\%$ difference between the $D^0 \to \rho^0\gamma$ and $D^0 \to \omega\gamma$ widths.

We estimate the long distance contribution $A(d\bar{d})$ from the calculation of the decays $D^0 \to \rho^0\gamma$, $\omega\gamma$ derived by the use of the HQET and chiral Lagrangian together with factorization in Refs. [7, 8]. The relative signs of different contributions to the parity conserving and parity violating amplitudes were left undetermined in Refs. [7, 8] and are taken from the quark models, as in [9], giving

$$A_{PC}(d\bar{d}) \simeq -3.7 \times 10^{-9}\ \mathrm{GeV}^{-1}\ , \qquad A_{PV}(d\bar{d}) \simeq 9.5 \times 10^{-9}\ \mathrm{GeV}^{-1}\ . \tag{15}$$

Inserting the numerical values of $A(d\bar{d})$ (15) and ϵ (6) to the expression for $\mathcal{D}^{\omega-\rho}$ (13) we get

$$\mathcal{D}^{\omega-\rho} \simeq \frac{0.025 + 0.54\ \mathrm{Re}\eta_{PC} + 3.6\ \mathrm{Re}\eta_{PV} - 0.003\ |\eta_{PC}|^2 - 0.02\ |\eta_{PV}|^2}{1 + 0.27\ \mathrm{Re}\eta_{PC} + 1.8\ \mathrm{Re}\eta_{PV} - 0.02\ \mathrm{Im}\eta_{PC} - 0.13\ \mathrm{Im}\eta_{PV} + 0.14\ |\eta_{PC}|^2 + 0.9\ |\eta_{PV}|^2} \tag{16}$$

where all orders in $\eta_{PC/PV}$ are retained. This expression displays the sensitivity of the quantity $\mathcal{D}^{\omega-\rho}$, defined by (13), on the magnitude of the short distance rate $c \to u\gamma$ contained in $\eta_{PC/PV}$ (10). We note that even if we neglect $\eta_{PC/PV}$, there is a difference of about 2.5% in the decay widths due to the $\rho - \omega$ mixing (2, 6), and another possible difference coming from $SU(3)$ breaking; together these amount to at most 15%.

The standard model prediction for the short distance contribution $A^{SD}(u\bar{u})$ is extracted from the calculation of the $c \to u\gamma$ rate at the two loop level [14] giving $Br(c \to u\gamma) \sim 3 \times 10^{-8}$ for D^0 decays. The matrix element $\langle \rho^0, \omega | \bar{u}\sigma_{\mu\nu}(1+\gamma_5)c | D^0 \rangle$ is calculated using the procedure described in detail in Ref. [7], giving

$$\begin{aligned} A^{SD}_{PC}(u\bar{u}) &\simeq (-1.4 \times 10^{-10} - i\ 4.0 \times 10^{-10})\ [1 \pm 0.2]\ \mathrm{GeV}^{-1}\ , \\ A^{SD}_{PV}(u\bar{u}) &\simeq (-2.3 \times 10^{-10} - i\ 6.6 \times 10^{-10})\ [1 \pm 0.2]\ \mathrm{GeV}^{-1}\ . \end{aligned} \tag{17}$$

The mechanism, shown in Fig . 1c, gives rise to a long distance part of $A(u\bar{u})$ and has been estimated in [7, 9]

$$A^{LD}_{PC}(u\bar{u}) \simeq (2.2 \pm 2.2) \times 10^{-10}\ \mathrm{GeV}^{-1}\ , \qquad A^{LD}_{PV}(u\bar{u}) \simeq (3.7 \pm 3.7) \times 10^{-10}\ \mathrm{GeV}^{-1}\ . \tag{18}$$

This contribution is small as it is proportional to the $SU(3)$ flavour breaking parameter

$$\frac{g_\rho^2(0)}{2m_\rho^2} - \frac{g_\omega^2(0)}{6m_\omega^2} - \frac{g_\phi^2(0)}{3m_\phi^2} \simeq (-1.2 \pm 1.2) \times 10^{-3}\ \mathrm{GeV}^2 \tag{19}$$

with g_V defined as $\langle V(q,\epsilon)|j_V^\mu|0\rangle = g_V(q^2)\epsilon^{*\mu}$. The mean value and error in (19) are determined from the experimental data on $V^0 \to e^+e^-$ decays [16] by assuming $g_V(0) \simeq g_V(m_V^2)$. Given that $A(u\bar{u}) = A^{SD}(u\bar{u}) + A^{LD}(u\bar{u})$ (17, 18), the standard model prediction for $\eta_{PC/PV}$ is

$$\begin{aligned} \text{Re } \eta_{PC} &\simeq -0.02 \pm 0.06 \ , \qquad \text{Im } \eta_{PC} \simeq 0.11 \pm 0.02 \\ \text{Re } \eta_{PV} &\simeq 0.01 \pm 0.04 \quad \ , \qquad \text{Im } \eta_{PV} \simeq -0.07 \pm 0.01 \end{aligned} \tag{20}$$

leading to the approximate equality

$$Br[D^0 \to \rho\gamma] \simeq Br[D^0 \to \omega\gamma] \simeq 1.2 \times 10^{-6} \ . \tag{21}$$

The standard model prediction for the relative difference of the rates $\mathcal{D}^{\omega-\rho}$ (13, 16) is

$$\mathcal{D}_{SM}^{\omega-\rho} \simeq 6 \pm 15 \ \ \% \ , \tag{22}$$

where the error is dominated by the uncertainty of the $SU(3)$ flavour breaking parameter, including (19). The strong rescattering, which can transform a $d\bar{d}$ pair to a $u\bar{u}$ pair does not change this result. This is because we consider decays into strong interaction eigenstates ρ and ω, which evolve as $\exp[-i(E - i\frac{1}{2}\Gamma)t]$.

For completeness, we consider also the case, when the ρ and ω states are experimentally identified only by the 2π and 3π decay modes, respectively, and not by their masses and widths. If the two pion final states, which arise from the decay of ω, are not disentangled in the experiment, then one is dealing with the final states of good isospin. A final state, which starts out as $\rho^{(I=1)}$, will obtain an admixture of $\omega^{(I=0)}$ due to the isospin mixing (2), and vice versa. The production of $2\pi\gamma$ and $3\pi\gamma$ final states depends on time. In order to extract the short distance contribution, we propose to look at the number of $2\pi\gamma$ and $3\pi\gamma$, integrated over time

$$\mathcal{D}^{3\pi-2\pi} \equiv \frac{\frac{N[D^0\to 3\pi\ \gamma]}{(m_D^2-m_\omega^2)^3 Br(\omega\to 3\pi)} - \frac{N[D^0\to 2\pi\ \gamma]}{(m_D^2-m_\rho^2)^3}}{\frac{N[D^0\to 3\pi\ \gamma]}{(m_D^2-m_\omega^2)^3 Br(\omega\to 3\pi)}} \ , \tag{23}$$

where 2π and 3π have invariant masses covered by ρ and ω resonances. At the first order in $\eta_{PC,PV}$ and ϵ we get in the standard model

$$\mathcal{D}_{SM}^{3\pi-2\pi} \simeq 4\ \frac{|A_{PC}(d\bar{d})|^2 \text{ Re } \eta_{PC} + |A_{PV}(d\bar{d})|^2 \text{ Re } \eta_{PV}}{|A_{PC}(d\bar{d})|^2 + |A_{PV}(d\bar{d})|^2} \simeq 4 \pm 15 \ \ \% \ , \tag{24}$$

where we have taken into account that $\epsilon(E_\omega - E_\rho)/\Gamma_\rho \ll \eta_{PC,PV}$.

The quantities $\mathcal{D}^{\omega-\rho}$ (13) and $\mathcal{D}^{3\pi-2\pi}$ (23) are particularly sensitive to new physics scenarios, which could enhance the $c \to u\gamma$ rate. Non-minimal realizations of the supersymmetric models, discussed in [2], can enhance these quantities up to

$$\mathcal{D}^{\omega-\rho} \simeq \mathcal{D}^{\omega^0-\rho^1} \simeq 1 \ . \tag{25}$$

Similar reasoning may be useful in extracting the FCNC transition $c \to ul^+l^-$ from the difference in the decay rates $\Gamma[D^0 \to \rho^0 l^+l^-]$ and $\Gamma[D^0 \to \omega l^+l^-]$, where l denotes an electron or muon. Long distance contributions to these have been studied in [9, 12] and arise via the mechanisms illustrated in Figs. 1a and 1c, where the photon is replaced with the virtual photon decaying to leptons. The corresponding differential branching ratio has resonant shape in terms of the di-lepton mass m_{ll} with maximums at $m_{ll} = m_\rho$, m_ω and m_ϕ. The LD contributions, which give rise to the final state $(d\bar{d})l^+l^-$, largely cancel in the difference $\Gamma[D^0 \to \omega^0 l^+ l^-] - \Gamma[D^0 \to \rho l^+ l^-]$, which is proportional to $A(u\bar{u}) = A^{SD}(u\bar{u}) + A^{LD}(u\bar{u})$. In this case, however, the $SU(3)$ flavour cancellation in $A^{LD}(u\bar{u})$ is not so effective as the maximums at $m_{ll} = m_{\rho,\omega}$ and $m_{ll} = m_\phi$ are well separated[2]. As a consequence, $|A^{LD}(u\bar{u})|$ is more than one order of magnitude larger than $|A^{SD}(u\bar{u})|$ and overshadows the interesting FCNC transition $c \to u l^+ l^-$ in the difference $\Gamma[D^0 \to \omega^0 l^+ l^-] - \Gamma[D^0 \to \rho l^+ l^-]$.

To summarize, we have proposed here yet another test for physics beyond the standard model in the charm sector. Our test exploits the extreme smallness of the $c \to u\gamma$ transition in standard model [14] and the near equality of the $D \to \rho^0\gamma$, $D^0 \to \omega\gamma$ amplitudes, obtained from the calculation [7, 8, 9] of long distance contributions to these decays. This equality would be spoiled, if, as encountered in certain supersymmetric models [2], the $c \to u\gamma$ amplitude is enhanced by a seizable factor. Standard model calculations, including $SU(3)$ breaking in the long distance contribution to the amplitudes, show that $D \to \rho^0\gamma$, $\omega\gamma$ widths do not differ by more than 15%. Thus, we are led to claim, conservatively, that a difference larger than 30% would be a "smoking gun" indication of new physics. The formalism leading to this conclusion has been exposed in this paper. In view of the expected $\sim 10^{-6}$ branching ratio for these decays, we look forward to the experimental tests, keeping in mind that the present upper limits are around 10^{-4} [18].

[2]The standard model predictions presented in [9, 12] give $Br^{LD}(u\bar{u}) \sim 10^{-7}$ and $Br^{SD}(u\bar{u}) \sim 10^{-9}$.

References

[1] G. Burdman, E. Golowich, J. L. Hewett and S. Pakvasa, Phys. Rev. D 52 (1995) 6383.

[2] I. Biggi, F. Gabbiani and A. Masiero, Z. Phys. C. **48** (1990) 633.

[3] S. Pakvasa, in Proc. FCNC 97 Conf. Santa Monica, Feb. 1997, World Scientific, 1998; hep-ph/9705397; J. L. Hewett, Proc. 28 Intern. Conf. on High Energy Physics (ICHEP 96), Warsaw, Poland, 1996, Z. Ajduk, A. K. Wroblewski (Eds.), World Scientific, 1997, Vol. 2, p. 1143.

[4] H. -Y. Cheng et al., Phys. Rev. D **51** (1995) 1199.

[5] A. Khodjamirian, G. Stoll and D. Wyler, Phys. Lett. B **358** (1995) 129.

[6] B. Bajc, S. Fajfer and R. J. Oakes, Phys. Rev. D **54** (1996) 5883.

[7] S. Fajfer and P. Singer, Phys. Rev. D **56** (1997) 4302.

[8] S. Fajfer, S. Prelovsek and P. Singer, Eur. Phys. J. C **6** (1999) 471; 715 (E).

[9] S. Prelovsek, Ph. D. thesis, to be published on `hep-ph`.

[10] K. S. Babu, X. G. He, X. -Q. Li and S. Pakvasa, Phys. Lett. B **205** (1988) 540.

[11] A. J. Schwartz, Mod. Phys. Lett. A **8** (1993) 967.

[12] S. Fajfer, S. Prelovsek and P. Singer, Phys. Rev. D **58** (1998) 094038.

[13] S. Fajfer, S. Prelovsek and P. Singer, Phys. Rev. D **59** (1999) 114003.

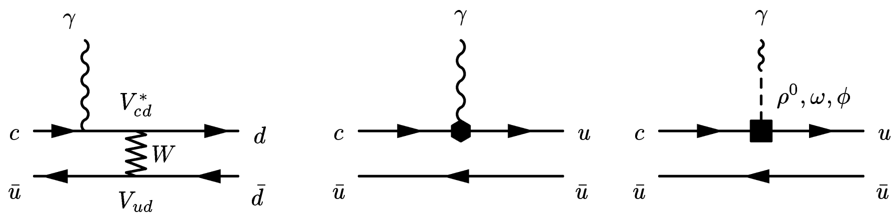
[14] G. Greub, T. Hurth, M. Misiak and D. Wyler, Phys. Lett. B **382** (1996) 415.

[15] Q. Ho -Kim and X. Y. Pham, Phys. Rev. D **61** (2000) 013008.

[16] Particle Data Group, C. Caso et. al, Europ. Phys. Journal C **3** (1998) 1.

[17] H. B. O'Connel *et al.*, Prog. Part. Nucl. Phys. 39 (1997) 201; Nucl. Phys. A 623 (1997) 559; K. Maltman *et al.*, Phys. Lett. B 376 (1996) 19; S. Gardner and H. B. O'Connel, Phys. Rev. D 59 (1999) 076002.

[18] D. M. Asner *et al.*, CLEO Collab., Phys. Rev. D 58 (1998) 092001.



(a) Long distance contribution leading to the $(d\bar{d})\gamma$ final state. The photon can be emitted from any of the quark lines.

(b) Short distance contribution arising from the FCN decay $c \to u\gamma$. The hexagon denotes the loop diagrams responsible for the $c \to u\gamma$ transition.

(c) Long distance contribution leading to the $(u\bar{u})\gamma$ final state. The box denotes the action of the weak nonleptonic effective Lagrangian.

Figure 1: Illustration of different contributios to the decays $D^0 \to \rho^0\gamma$ and $D^0 \to \omega\gamma$.