Z. Phys. C - Particles and Fields 16, 279-281 (1983)

Zeitschrift für Physik C Particles and Fields

On the Radiative Decay of Orthoquarkonium Via Two Intermediate Gluons: $1^{--}(Q\bar{Q}) \to \gamma + 1^{++}(q\bar{q})$

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Received 28 September 1982

Abstract. We investigate the decay of a heavy $^3S_1(Q\bar{Q})$ vector meson into a real photon and two off-shell intermediate gluons which create a $^3P_1(q\bar{q})$ axial vector meson. As an application we suggest to look for the decay $J/\psi \to \gamma + D(1285)$ and $\Upsilon \to \gamma_1 + P_c/\chi_{1^{++}} \to \gamma_1 + \gamma_2 + J/\psi$.

The transitions $e^+e^- \to \gamma^* \to J/\psi \to \gamma$ (direct) + meson are an 'ideal laboratory' to study the properties of $C = +$ and $I = 0$ mesons. Also, since the J/ψ is polarized due to its formation from e^+e^- the polar angular distribution of the γ (or meson) yields already some information on the spin of the meson:

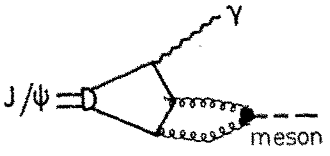
$$J_{\text{meson}} \geqq 1 \quad \text{if } W(\Theta\gamma e^{\pm}) \neq 1 + \cos^2\Theta\gamma e^{\pm}.$$

Up to now already several mesons have been clearly identified in the radiative J/ψ decays:

$$J^P = 0^-: \eta, \eta', \mathscr{I}(1440); \qquad J^P = 2^+: f, \Theta(1640) \text{ [1]}$$

with branching ratios $B(J/\psi \to \gamma + \text{meson})$ on the order of 10^{-3}. With present or future high statistics experiments it should be possible to identify meson states occurring with branching ratios of the order 10^{-4}.

In lowest order QCD these Zweig forbidden transitions are mediated by two gluon exchange



+permutations (1)

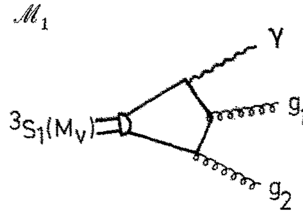
where the coupling of the gluons is either to 'light quarkonia' or 'gluonia'.

In this paper we present the results for the transition rate and helicity structure of

$$\begin{aligned} &^3S_1(Q\bar{Q}) \to \gamma + {}^3P_1(q\bar{q}) \\ &1^{--}(Q\bar{Q}) \to \gamma + 1^{++}(q\bar{q}) \end{aligned} \tag{2}$$

by using (1) as the transition mechanism. The Landau-Pomeranchuk-Yang theorem[2] tells us that the transition is mediated only if at least one of the gluons is off-shell. Thus a 1^{++} state does not occur in the $1^{--} \to \gamma + g_1 + g_2$ Born term: $k_1^2 = k_2^2 = 0$. But the Feynman rules in (1) yield a nonvanishing dispersive part stemming from the box integration. That the off-shell part of the loop contribution can be appreciable or even dominating has already been demonstrated in the transition $^3S_1(Q\bar{Q}) \to \gamma + {}^1S_0(q\bar{q})$ using again (1) as dynamical input[3].

The calculation of the decay amplitude (1) for a 1^{++} meson in the final state is performed by contracting the (off-shell) Ore-Powell matrix element $\mathscr{M}_1$



$$\begin{aligned} \mathscr{M}_1 = &\frac{i8\sqrt{M_v}\cdot\psi(0)}{(k_1+k_2,k_3)(k_2+k_3,k_1)(k_3+k_1,k_2)} \\ &\cdot \{(\varepsilon_1^*\varepsilon_2^*)[-(k_1k_3)(\varepsilon_3^*k_2)(\varepsilon_v k_1) - (k_1k_3)(k_2k_3)(\varepsilon_v\varepsilon_3^*) \\ &-(k_2k_3)(\varepsilon_3^*k_1)(\varepsilon_v k_2)] + (\varepsilon_v\varepsilon_3^*)[(k_3k_1)(\varepsilon_1^*k_2)(\varepsilon_2^*k_3) \\ &-(k_1k_2)(\varepsilon_1^*k_3)(\varepsilon_2^*k_3) + (k_3k_2)(\varepsilon_2^*k_1)(\varepsilon_1^*k_3)] \\ &+1 \leftrightarrow 3 + 2 \leftrightarrow 3\} \end{aligned} \tag{3}$$

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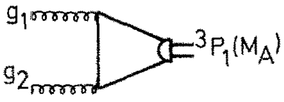
with the 1^{++} decay element



$$\mathcal{M}_2^* = \frac{-i2\sqrt{6M_A}\cdot\psi_P'(0)}{M_A^3\cdot(k_1k_2)^2}$$
$$\cdot\{(k_1k_2)(k_1^2-k_2^2)\epsilon(\varepsilon_A^*,\varepsilon_1,\varepsilon_2,P_A)+[2(k_1k_2)(k_2\varepsilon_2)$$
$$-(k_1^2+k_2^2)(k_1\varepsilon_2)]\epsilon(\varepsilon_A^*,\varepsilon_1,P_A,k_1-k_2)+1\leftrightarrow 2\} \tag{4}$$

leading to the transition amplitude

$$\mathcal{T}=\frac{-1}{(2\pi)^4}\int d^4k_1d^4k_2\,\delta(P_A-k_1-k_2)\frac{\mathcal{M}_1\cdot\mathcal{M}_2^*}{k_1^2k_2^2} \tag{5}$$

A straightforward evaluation of (5) leads to a form involving at most three parameter Feynman integrals. The results agree with an alternative evaluation using covariant helicity projectors which reduce the number of necessary Feynman parameter integrations to two. The two independent helicity amplitudes labelled by the helicity of the 1^{++} particle

$$1^{--}(1,0)\to\gamma(1)+1^{++}(0,1)\equiv(\mathcal{H}_0,\mathcal{H}_1) \tag{6}$$

are given by

$$\begin{pmatrix}\mathcal{H}_0\\ \mathcal{H}_1\end{pmatrix}=\mathcal{N}\begin{pmatrix}\mathcal{H}_0(x)\\ \mathcal{H}_1(x)\end{pmatrix}=\mathcal{N}\begin{pmatrix}\sqrt{1-x}\,A_1(x)\\ \frac{2-x}{2}A_1(x)-\frac{x}{2}A_2(x)\end{pmatrix}$$

$$x=1-M_A^2/M_v^2$$

$$\mathcal{N}=8\sqrt{M_v}\,\psi(0)\cdot\frac{2\sqrt{6M_A}}{M_A^3}\psi_P'(0)\cdot\frac{16\pi^2}{(2\pi)^4M_v}$$
$$\cdot\, e_Q(4\pi)^{\frac{5}{2}}\sqrt{\alpha}\cdot\alpha_s(M_v)\cdot\alpha_s(M_A)\cdot\tfrac{2}{3}\cdot\tfrac{1}{2} \tag{7}$$

and where

$$A_1(x)=\frac{2(1-x)^2}{x^3}\log(1-x)-\frac{2(1-x)}{x^2}\log(1-x)$$
$$+\frac{2(1-x)}{x^2}\log 2x+\frac{3x-2}{x(1-2x)}\log 2x$$
$$+\frac{1}{x^3}[Li_2(1)-Li_2(1-2x)]$$
$$+\frac{6(1-x)}{x^2}[Li_2(1-2x)-Li_2(1-x)+\log 2\log(1-x)]$$

$$A_2(x)=\frac{-8(1-x)^2}{x^3}\log(1-x)+\frac{2(1-x)}{x^2}$$
$$-\frac{7(1-x)}{x^2}\log 2x-\left(1-\frac{5}{x}+\frac{3}{x^2}\right)\cdot\frac{\log 2x}{1-2x}$$
$$+\frac{1}{x^3}[Li_2(1)-Li_2(1-2x)]+(1-x)\cdot\left(\frac{12}{x^3}-\frac{10}{x^2}\right)$$
$$\cdot\,[Li_2(1-2x)-Li_2(1-x)+\log 2\log(1-x)]$$
$$\left(Li_2(0)=0,\ Li_2(1)=\frac{\pi^2}{6}\right)$$

The limiting behaviour of $A_{1,2}$ is given by

$$\lim_{x\to 1}A_1=\lim_{x\to 1}A_2=\frac{\pi^2}{4}-\log 2;$$
$$\lim_{x\to 0}A_1=\lim_{x\to 0}A_2=-\tfrac{2}{3}\log 2x+\tfrac{13}{18}$$
$$A_1(\tfrac{1}{2})=1+\frac{\pi^2}{3}-6\log^2 2;$$
$$A_2(\tfrac{1}{2})=7-\pi^2+16\log 2-14\log^2 2. \tag{8}$$

We observe that in the limit $M_A\to 0$ the amplitude $\mathcal{H}_0\propto M_A^{-3/2}$, $\mathcal{H}_1\propto M_A^{-1/2}$, i.e. the longitudinal contribution dominates in this limit as expected. The ratio H_1/H_0 is shown in Fig. 1. The angular distribution of the γ or 1^{++} with respect to the electron-positron beam in $e^+e^-\to Q\bar{Q}\to\gamma+1^{++}$ is $W(\Theta)\propto 1+\alpha(x)\cos^2\Theta$ where $\alpha(x)$ is given by $(1-2H_1^2/H_0^2)/(1+2H_1^2/H_0^2)$. For small x it approaches the electric dipole limit. Figure 2 gives the combination $x(H_0^2$

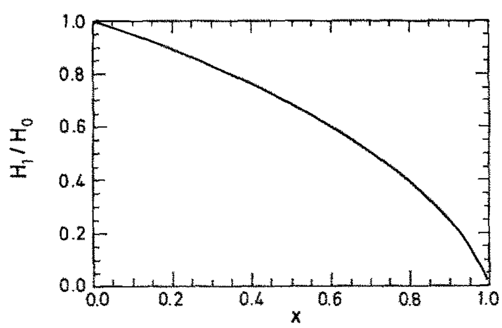


Fig. 1. See text

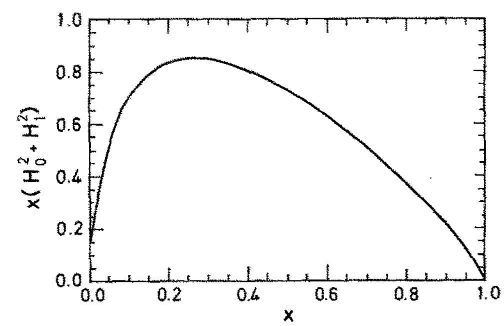


Fig. 2. See text

$+H_1^2$) which occurs in the rate formula

$$\Gamma=\frac{1}{24\pi}\cdot\frac{x}{M_v}(\mathscr{H}_0^2+\mathscr{H}_1^2) \tag{9}$$

We have estimated the rate $J/\psi\to\gamma+D(1285)$ to be of the order of 16 eV using $\alpha_s(M_{J/\psi})=0.2$ and $\alpha_s(M_D)=0.3$ and extracting $\frac{|\psi_P'(0)|^2}{M_D^4}=0.25\,\mathrm{MeV}$ from $\Gamma(f\to\gamma\gamma)=3\,\mathrm{keV}$.

Acknowledgements. We thank R. Kleiss and H. Schneider for some helpful discussions.

Note Added in Proof. If experimentally $D(1285)$ is not found at this level it indicates, in our opinion, a different wave function for the 1^{++} than the $^3P_1$ quark-antiquark wave function used here. Since the D is in the same $SU(3)$ multiplet as the A_1 and the A_1 is known from the τ-lepton decay to couple to the axial vector current $\gamma_5\gamma_\mu$, a relativistic wave function for the 1^{++} of the type $\gamma_5\not{\varepsilon}_A$ may be more realistic. However, in such a case a 1^{++}-state would not couple to two gluons since trace $\{\gamma_5\not{\varepsilon}_A\not{\varepsilon}_1\not{\varepsilon}_2\}=0$. A candidate to check our nonrelativistic calculation would then be the transition $\gamma\to\gamma_1+P_c/\chi_{1^{++}}$ where in addition to γ_1, the γ_2 from the $P_c/\chi_{1^{++}}\to\gamma_2+J/$ cascade decay can be used as a good trigger.

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