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Decays of W and Z in the broken-colour model

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Abstract. The two-gluonic decay modes of W and Z in the broken colour model with integrally charged quarks are considered. The gluonic branching ratios are found to be 3 % and 2·7 % for W and Z respectively. The angular distributions of the decays of W and Z to two jets of hadrons are also worked out.

Keywords. Integer-charge quarks; broken colour; weak bosons.

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The experimental detection of the W and Z bosons has already been reported (UA1 Collaboration 1983a, b; UA2 Collaboration 1983a, b). Hopefully, detailed experimental data on the various decay modes of these electroweak gauge bosons will become available soon. The theoretical expectations for these decay modes in the framework of the standard $SU(3) \times SU(2) \times U(1)$ model with unbroken colour have already been worked out (Quigg 1977; Kajantie 1982). In this note, we look at these decay modes of the electroweak gauge bosons following from the $SU(3) \times SU(2) \times U(1)$ model with broken colour, colour being broken in such a way that the quarks are integrally charged.

The spontaneous breaking of colour through the Higgs mechanism was discussed and the consequences for deep inelastic lepton-hadron scattering as well as e^+e^- annihilation were analysed in the early papers (Pati and Salam 1973, 1976; Rajasekaran and Roy 1975, 1976). Subsequently, the broken-colour model was confronted with two-jet and three-jet phenomena in e^+e^- annihilation and it was shown that the model cannot be ruled out by present data (Rajasekaran and Rindani 1979, 1980, 1981, 1982; Lakshmibala *et al* 1981). It has been shown that even the two-photon experiments are not decisive in this respect because of the virtualness of the photon (Jayaraman *et al* 1982). In fact, the excess of events in the PLUTO data on two photon experiments can be interpreted (Godbole *et al* 1984) as positive evidence in favour of the broken-colour model. In view of these, it may be reasonable to regard the broken-colour model as an alternative to standard QCD at least at the phenomenological level, until the latter is established to be *the* correct theory.

While the standard model with unbroken colour and fractionally charged quarks allows in the hadronic sector only decays like $Z \to q\bar{q}$ and $W \to q\bar{q}$ to order G_F and zeroth order in the strong coupling constant, the broken-colour model with integrally charged quarks allows to the same order, the decays $Z \to gg$ and $W \to gg$, where g refers to the gluon. This is because of the mixing between the gluons and the electroweak gauge bosons, which is a characteristic feature of the broken colour theories.

Although there exists an additional piece in the $q\bar{q}W$ (and $q\bar{q}Z$) vertex in the broken-colour model arising from the quark current, this is suppressed for $q^2 \gg m_g^2$, where q is the momentum of the current and m_g is the gluon-mass parameter. Hence for $m_W^2 \gg m_g^2$ and $m_Z^2 \gg m_g^2$, this colour octet vertex can be ignored and the widths $\Gamma(W/Z \to q\bar{q})$ are given by the colour singlet vertex alone and therefore these decays are the same as in the standard model.

The only new effect in the broken-colour model is therefore the contribution coming from the gluons mentioned above. In this note, this gluonic contribution to the two-jet decay modes of W and Z is presented. This gluonic contribution is actually independent of the details of the specific mechanism of breaking of colour-symmetry in the integer-charge quark model.

It is worth pointing out that this gluonic contribution exists only above the colour threshold. Because of the possibility of the energies of the order of 80 GeV being above the colour threshold, the additional contribution to the W and Z decays may provide a new window on colour to be explored.

Let θ be the angle between the direction of polarisation of the electroweak boson and the direction of emission of either gluon, in the rest frame of the electroweak boson. Then the decay angular distribution for the gluon emission from the electroweak boson is given by

$$\frac{d\Gamma}{d\cos\theta}(B \to gg) = \frac{1}{4\pi}\frac{C_B^2}{8m_B^2}\left(\frac{m_g^2}{m_B^2 - m_g^2}\right)^2\left(\frac{m_B^2}{4} - m_g^2\right)^{3/2} \times \left[\frac{8m_B^2}{m_g^2} + \left(\frac{m_B^4}{2m_g^4} - \frac{2m_B^2}{m_g^2} + 6\right)\sin^2\theta\right], \tag{1}$$

where B refers to the electroweak boson W or Z, m_B and m_g are respectively the masses of the electroweak boson and gluon and C_B is the Bgg coupling constant in the broken-colour model. In terms of the semiweak coupling constant g and the electroweak mixing angle θ_W, C_B is given by (Rajasekaran and Roy 1975).

$$C_W = \sqrt{3}g, \tag{2a}$$

and

$$C_Z = g\frac{(3 - 6x_W + 4x_W^2)^{1/2}}{(1 - x_W)^{1/2}}, \tag{2b}$$

where $x_W = \sin^2\theta_W$.

The angular distribution in (1) does not contain terms odd in $\cos\theta_w$ and so there is no up-down asymmetry. This is due to the fact that the Bgg coupling is parity conserving.

It should be noted that in the broken-colour theory, the gluons do have mass m_g, but this is an effective mass parameter. Taking it to be small (as indicated by the analysis (Jayaraman *et al* 1982; Godbole *et al* 1984) of two-photon data), and ignoring it in comparison to m_B we get

$$\frac{d\Gamma}{d\cos\theta}(B \to gg) = \frac{1}{4\pi}\frac{C_B^2 m_B}{128}\sin^2\theta, \tag{3}$$

and integrating over the angle,

$$\Gamma(B \to gg) = \frac{1}{4\pi}\frac{C_B^2 m_B}{96}. \tag{4}$$

The factor $[m_g^2/(m_B^2-m_g^2)]^2$ in (1) arises from the mixing between the gluons and the electroweak bosons already referred to and it can be seen that this factor is essential for the m_g-independent results given in (3) and (4).

In terms of G_F, we have

$$\Gamma(W\to gg)=\frac{G_F}{\sqrt{2}}\frac{m_W^3}{16\pi}\approx 82{\cdot}4\,\text{MeV}, \tag{5}$$

$$\Gamma(Z\to gg)=\frac{G_F}{\sqrt{2}}(3-6x_W+4x_W^2)\frac{m_Z^3}{48\pi}\approx 74{\cdot}8\,\text{MeV}, \tag{6}$$

where we have used

$$x_W=0{\cdot}22,\ m_W=79{\cdot}5\,\text{GeV}\quad\text{and}\quad m_Z=90\,\text{GeV}.$$

These are the additional contributions to the widths. The gluonic branching ratios are*

$$B(W\to gg)=\frac{1}{\frac{8}{3}(N_l+3N_q)+1}=\frac{1}{33}, \tag{7}$$

and

$$B(Z\to gg)=\frac{c_W}{4a_WN_l+12b_WN_q+c_W}\approx\frac{1}{36{\cdot}6}, \tag{8}$$

where N_l and N_q denote the number of lepton doublets and quark doublets respectively (taken here to be three each) into which the weak boson can decay and a_W, b_W and c_W are known functions of x_W given by

$$a_W=2-4x_W+8x_W^2, \tag{9a}$$

$$b_W=2-4x_W+\frac{40}{9}x_W^2, \tag{9b}$$

and

$$c_W=3-6x_W+4x_W^2. \tag{9c}$$

These branching ratios are comparable to the branching ratios for decay into certain leptonic modes such as $Z\to e^+e^-$ and hence not negligible.

For the polarised W, the angular distribution for the gg mode is given by (3) as

$$\frac{\mathrm{d}\Gamma}{\mathrm{d}\cos\theta}(W\to gg)=\frac{3}{64\pi}\frac{G_F}{\sqrt{2}}m_W^3(1-\cos^2\theta). \tag{10}$$

This can be compared with the angular distribution for the $q\bar{q}$ mode (Quigg 1977; Kajantie 1982)

$$\frac{\mathrm{d}\Gamma}{\mathrm{d}\cos\theta}(W\to q\bar{q})=\frac{9}{16\pi}\frac{G_F}{\sqrt{2}}m_W^3(1-\cos\theta)^2, \tag{11}$$

* We have ignored the masses of all the leptons and quarks in comparison with those of the electroweak bosons. For a massive top-quark, with $m_t=30$ GeV the numbers on the right side of (7) and (8) will be 1/31·3 and 1/34·4 respectively. In the extreme case of a much heavier top-quark, the top-quark mode can be omitted altogether, in which case these numbers will become 1/25 and 1/32·5 respectively.

where the number of quark-doublets N_q has been taken to be 3 and quark masses have been ignored. Here θ is the angle between the polarisation direction of the W and the direction of emission of the quark. Equation (11) has a parity violating $\cos\theta$ term. However if we consider jet production without distinguishing the quark from the antiquark this term will vanish (Schiller 1979) and we will have

$$\frac{d\Gamma}{d\cos\theta}(W\rightarrow \text{quark jets}) = \frac{9}{8\pi}\frac{G_F}{\sqrt{2}}m_W^3(1+\cos^2\theta). \tag{12}$$

Just like the quarks, the gluons are also expected to materialise in the form of hadronic jets and we will have

$$\frac{d\Gamma}{d\cos\theta}(W\rightarrow \text{gluon jets}) = \frac{3}{32\pi}\frac{G_F}{\sqrt{2}}m_W^3(1-\cos^2\theta). \tag{13}$$

Thus in the ICQ model the weighted angular distribution for $W\rightarrow 2$ hadronic jets is obtained by adding (12) and (13) to give

$$\frac{d\Gamma}{d\cos\theta}(W\rightarrow 2\text{ jets}) = \frac{39}{32\pi}\frac{G_F}{\sqrt{2}}m_W^3(1+\tfrac{11}{13}\cos^2\theta). \tag{14}$$

The corresponding distributions for polarised Z are

$$\frac{d\Gamma}{d\cos\theta}(Z\rightarrow gg) = \frac{1\cdot 87}{64\pi}\frac{G_F}{\sqrt{2}}m_Z^3(1-\cos^2\theta), \tag{15}$$

and

$$\frac{d\Gamma}{d\cos\theta}(Z\rightarrow q\bar{q}) = \frac{1\cdot 5}{4\pi}\frac{G_F}{\sqrt{2}}m_Z^3(1-1\cdot 67\cos\theta+\cos^2\theta). \tag{16}$$

As in the case of the W boson, we can get the weighted angular distribution for $Z\rightarrow 2$ jets of hadrons to be

$$\frac{d\Gamma}{d\cos\theta}(Z\rightarrow 2\text{ jets}) = \frac{1\cdot 62}{2\pi}\frac{G_F}{\sqrt{2}}m_Z^3(1+0\cdot 86\cos^2\theta). \tag{17}$$

In principle these jet angular distributions given here may serve to distinguish the two models. However, it should be remarked that just as in the other calculations to zeroth order in the strong coupling in broken-colour theory reported so far (Pati and Salam 1976; Rajasekaran and Roy 1975, 1976; Rajasekaran and Rindani 1979) these gluonic signals are rather weak and hence we conclude that even in the $W-Z$ sector, broken colour does not lead to any drastic difference at the phenomenological level.

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