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NEW MESONIC STATES NEAR 160 GeV

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A B S T R A C T

Recently reported data on monojet events and on massive $e\nu$ jet and jet-jet systems at the CERN SPS $\bar{p}p$ collider suggest a common origin which could be the production of new hadronic states with mass in the neighbourhood of 160 GeV. We explore the consequences of postulating that such states exist and are excited in $q\bar{q}$ or gg subprocesses. To be specific, we identify the states as coloured mesonic systems and we then compute production cross-sections and branching fractions into $q\bar{q}$, gg , gZ^0 , gW , and $g\gamma$, obtaining fair consistency with data. Implications and predictions are discussed.

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Events with very interesting new topologies have been found recently in data from the CERN UA1¹⁾ and UA2²⁾ collaborations. Hadronic jets are identified along with substantial missing transverse energy¹⁾, with a single additional lepton prominent in some cases²⁾. In addition, at least two other phenomena are noteworthy. Evidence has been presented²⁾ for an enhancement in the jet-jet invariant mass distribution in the vicinity of 150 GeV. Second, there may be differences in the Z^0 and $W^\pm$ production mechanisms, reflected either as a tendency for the ratio of Z^0 and W yields to be somewhat larger than expected³⁾, or in extra recoil jet activity in Z^0 events. Such differences are not expected from the prototype Drell-Yan production mechanism.

In theoretical interpretations proposed so far for some of these phenomena, at least one major assumption is invoked, e.g., supersymmetry⁴⁾, excited quarks⁵⁾, enhanced Higgs couplings⁶⁾, bound states of vector bosons⁷⁾. In this letter, we propose and discuss the possibility that features of the data mentioned above have a common origin in the production of massive mesonic states, with mass in the neighbourhood of 160 GeV. The data are consistent with the hypothesis that these new states may be colour octet mesonic systems, electrically charged as well as neutral, formed from fusion of the light quarks via the mechanism $q\bar{q} \rightarrow M_c$. The fusion of two gluons may also produce isosinglet states, but, at $\sqrt{s} = 540$ GeV and $M \approx 160$ GeV, the relevant luminosity is decidedly smaller than for $q\bar{q}$.

While the assumption of coloured mesons of such high mass may seem unusual, the data themselves are peculiar. The exploration of various alternative explanations, including coloured resonances, seems useful from the points of view of trying to correlate the diverse observations and of making predictions for other phenomena which may be studied in the current data sample and in future collider experiments at CERN and at Fermilab. For estimates of production rates and branching fractions, we take the new state M_c to be a pseudoscalar, but other spins and parities are possible, as well as a family thereof. Some of the decay modes, branching fractions and observational consequences are listed in the Table. The expected event rates are not inconsistent with current observations, and some tests can be made (e.g., γ jet events) even at the present luminosity level.

Our principal reason for postulating a coloured state is a desire to correlate the diverse data mentioned above, including the UA1 monojet events¹⁾ (here taken to arise from $M_c \rightarrow gZ^0, Z^0 \rightarrow \nu\bar{\nu}$) and the UA2 jet-lepton-missing transverse energy events²⁾ (here $gW^\pm$). Within the theoretical framework of QCD there is no reason for the existence of heavy coloured states, but no reason either to exclude them^{1*)}. Indeed, one may entertain models in which all colour singlet states have negligible mass while coloured states (colour-octet meson states in particular) are very massive. To illustrate this point, a simple model may be considered in which the mass is obtained as a sum of two contributions: a two-body potential term (with couplings proportional to those associated with one gluon, or, more generally, colour octet exchange) and an effective mass term proportional to the number of quarks. With the colour singlet $q\bar{q}$ and $q\bar{q}q$ states both set to practically zero mass, one would then obtain colour octet $q\bar{q}$ states at a mass expressed as $3m$, where m has an a priori large value. Additional consequences would be the existence of coloured diquark ($\bar{3}$) states with mass $\sim 4m/3$ and (6) states with mass $10m/3$. We here envision the colour octet states as extra pole singularities in the short range $q\bar{q}$ scattering amplitude, to be added as relatively important terms beyond those associated with lowest order QCD perturbation theory.

Once colour octet mesons (M_c) are postulated, several implications follow immediately. For example, monophoton events are not expected at an appreciable rate since colour forbids $M_c \rightarrow \gamma Z^0$. However γ -jet events are expected; these may only very rarely be observed as isolated photons. Likewise, M_c does not have decay modes such as $W^\pm W^\mp$, $Z^0 Z^0$, or $W^\pm Z^0$, which would be typical of, say, Higgs decay, nor are decay modes containing heavy flavours favoured. The $\bar{3}$ diquark states mentioned above would be expected to provide anomalies at a mass scale in the neighbourhood of 70 GeV, at a rate too low perhaps to be discernable above the large QCD "background" in the dijet invariant mass distribution in $p\bar{p}$ collisions (the more so for a very massive sextet coloured diquark). Other implications are mentioned later.

Our choice for the mass of M_c of the order of 160 GeV is based on partial evidence from the UA2 collaboration²⁾ of a peak in the jet-jet invariant mass distribution in this neighbourhood and the reported invariant masses of their

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Although we are working within the framework of confined colour, one should be open to new possibilities. Indeed, quark search experiments still take place!

jet-lepton-missing transverse energy events²⁾. The kinematics of the unusual UAL events¹⁾ are also compatible with this mass.

At $M = 160$ GeV and $\sqrt{s} = 540$ GeV ($\tau^{\frac{1}{2}} = M/\sqrt{s} \approx 0.30$), the differential luminosity for $u\bar{u}$ subprocesses in a $p\bar{p}$ collision is about a factor of 6 greater than that for $g\bar{g}$ subprocesses. It is natural therefore to seek an interpretation of production in which $q\bar{q}$ subprocesses play an important role, analogous to the production of colour singlet states via $q\bar{q}$ fusion, while $g\bar{g}$ subprocesses could, to a lesser extent, also feed isoscalar states. The colour factors, averaged over initial colours, are $4/9$ for $q\bar{q} \rightarrow M_c$ and $3/8$ for $g\bar{g} \rightarrow M_c$, a slight further advantage for production via $q\bar{q}$. According to the model discussed later, the squared matrix elements for $g\bar{g} \rightarrow M_c$ and $q\bar{q} \rightarrow M_c$ are about equal [c.f., Eqs. (14) and (18)]. The large mass of M_c produced through a large momentum transfer between quarks or gluons is assumed to justify a perturbative treatment of production, and large momentum transfers do the same for decay processes.

Standard arguments lead to the relationship

$$\sigma(p\bar{p} \rightarrow M_c X) = \tau \frac{dL(q\bar{q})}{d\tau} \hat{\sigma}(q\bar{q} \rightarrow M_c). \quad (1)$$

Here only the parton subprocess $q\bar{q} \rightarrow M_c$ has been retained, and $\tau(dL/d\tau)$ is the relevant differential luminosity. The production rate of events is then $N(p\bar{p} \rightarrow M_c X) = L^{p\bar{p}} \sigma(p\bar{p} \rightarrow M_c X)$, where $L^{p\bar{p}}$ is the luminosity of $p\bar{p}$ collisions.

We may further express $\hat{\sigma}$ in terms of the partial width $\Gamma(M_c \rightarrow q\bar{q})$, as follows:

$$\hat{\sigma} = \frac{2}{9} \frac{16\pi^2 \Gamma(M_c \rightarrow q\bar{q})}{M_c^3}; \quad (2)$$

M is the mass of the state M_c . The factor $2/9$ is a combination of spin and colour factors. If M_c were instead a colour singlet, the right-hand side of Eq. (2) would be reduced by a factor 8, and if M_c were assigned spin 1 instead of our assumed spin zero, there would be an increase by a factor 3. For comparison,

$$\hat{\sigma}(u\bar{d} \rightarrow W^+) = \frac{1}{3} \frac{4\pi^2 \Gamma(W^+ \rightarrow u\bar{d})}{M_W^3}. \quad (3)$$

Luminosity calculations at $\sqrt{s} = 540$ GeV show that $\hat{s}^{-1} \tau(dL/d\tau)$ drops by one order of magnitude from $\sqrt{\hat{s}} = M_W$ to $\sqrt{\hat{s}} = 160$ GeV. Combining Eqs. (1)-(3), we estimate the ratio (for $p\bar{p}$ at $\sqrt{s} = 540$ GeV):

$$\begin{aligned} \frac{\sigma(p\bar{p} \rightarrow M_c X)}{\sigma(p\bar{p} \rightarrow W^+ X)} &\approx \frac{8}{36} \frac{M_W}{M_c} \frac{\Gamma(M_c \rightarrow u\bar{d})}{\Gamma(W^+ \rightarrow u\bar{d})} \\ &\approx \frac{\Gamma(M_c \rightarrow u\bar{d})}{6}, \quad \text{in GeV units} \end{aligned} \quad (4)$$

To obtain the second expression in Eq. (4), we have used the standard expectation that $\Gamma(W \rightarrow u\bar{d}) \approx 780$ MeV.

General arguments suggest that for typical hadronic states both $\Gamma(M_c \rightarrow q\bar{q})$ and $\Gamma(M_c \rightarrow g\bar{g})$ are equal to $\alpha_s^2 M_c$ up to factors of order unity, whereas $\Gamma(M_c \rightarrow g\gamma) \sim \Gamma(M_c \rightarrow gZ^0) \sim \Gamma(M_c \rightarrow gW) \sim \Gamma(M_c \rightarrow gW) \sim \Gamma(M_c \rightarrow gW)$ are of order $\alpha_s \alpha_{em} M_c$. Therefore, the width of the state M_c should be of the order of a few GeV and, at collider energy, it would decay within the confinement range. Since about 600 W 's have been produced in each experiment (taken into account is the branching ratio of the leptonic decay modes seen so far), Eq. (4) shows that several hundred M_c 's could also have been formed, and highly visible decay modes at the percent level could be responsible for the unusual events reported so far. In the UA2 data²⁾, the width of the peak reported in the jet-jet mass distribution is consistent with the experimental resolution, implying $\Gamma_{\text{total}} \approx 13$ GeV.

Although obtained by simple arguments, Eq. (4) is quite important and is a major motivation for proposing a coloured state. It may be contrasted with the result applicable for a colour neutral heavy onium state excited by $g\bar{g}$

subprocesses. In this case, the right-hand side of Eq. (4) is replaced by $(2J+1)\Gamma(gg)/100$ in GeV units, where J is the spin of the onium state. This means that to obtain a production rate comparable to that of the W^+ , several onium states with large width would be required instead of our single colour octet meson state excited by qq subprocesses.

Lacking a detailed dynamical model for the structure of M_c , or even a phenomenological Lagrangian to provide its couplings to quarks and to gauge bosons, we have examined two different approaches to obtain phenomenological estimates of the partial decay widths. First, we considered a simple "spectator" model for $M_c \rightarrow W X$ and $M_c \rightarrow Z^0 X$, illustrated in Fig. 1, in which each quark constituent is imagined to decay independently. We ignore quark masses. Because the invariant matrix element squared and summed over spins is identical for $q_i \rightarrow W q_f$ and for $W \rightarrow ev$, we may conclude directly that (for each charge state of M_c)

$$\frac{\Gamma(M_c \rightarrow W^+ X)}{\Gamma(W^+ \rightarrow e^+ \nu)} = \frac{3}{2} \frac{M_W}{M} \frac{M^2 - M_W^2}{M^2} \approx 0.6. \quad (5)$$

The factor of 3/2 arises from the averages over spin states; the ratio M_W/M comes from flux factors in the computation of the Γ 's; and the third factor arises from the phase space integral for $M_c \rightarrow W X$ (the system X is assigned zero invariant mass for the sake of this discussion).

Using $\Gamma(W \rightarrow ev) = 260$ MeV provided by the standard model, we deduce

$$\Gamma(M_c \rightarrow W^+ X) \approx 160 \text{ MeV}; \quad (6)$$

and, with $\alpha_{em}/\alpha_s \approx 0.1$, we obtain

$$\Gamma(M_c \rightarrow g X) \approx 2 \text{ GeV}, \quad (7)$$

with $\Gamma(M_c \rightarrow q X)$ of about the same size. Reference to Eq. (4) then allows us to conclude that the production rates for M_c and W are within a factor of ~ 3 of being equal at $\sqrt{s} = 540$ GeV. In this connection, we may note that for the dominant jet-jet decay modes of W and M_c , the UA2 results²⁾ lead one to compare 69 ± 17 events in the 160 GeV "peak" to 110^{+81}_{-70} events which would be attributed to W and Z formation and decay. These two values are of the same order of magnitude.

When Eqs. (4) and (5) are combined, one may deduce that there should be about one event of the type $M_c \rightarrow W \text{ jet}$, $W \rightarrow ev$, for every 13 in which the W is produced directly (via the Drell-Yan mechanism), with subsequent $W \rightarrow ev$. This expected ratio is consistent with the observations of the UA2 collaboration^{2),3)} (32 ± 6 W $\rightarrow ev$ events, and three events which may be interpreted as $M_c \rightarrow W \text{ jet}$, $W \rightarrow ev$).

In a similar fashion, we may employ the spectator model to derive

$$\frac{\Gamma(M_c \rightarrow Z X)}{\Gamma(Z \rightarrow e^+ e^-)} \approx 1.25; \quad (8)$$

and, with $\Gamma(Z \rightarrow e^+ e^-) \approx 90$ MeV,

$$\Gamma(M_c \rightarrow Z X) \approx 100 \text{ MeV}. \quad (9)$$

Taken together, Eqs. (5) and (8) imply

$$\frac{\Gamma(M_c \rightarrow Z X)}{\Gamma(M_c \rightarrow W^+ X)} \approx 2/3, \quad (10a)$$

and

$$\frac{\Gamma(M_c \rightarrow ZX)}{\Gamma(W^+ \rightarrow e^+ \nu)} \simeq \frac{4}{9} \quad (10b)$$

The ratio in Eq. (10a) is in agreement with our more detailed but model dependent computation of $\Gamma(M_c \rightarrow Zg)/\Gamma(M_c \rightarrow Wg)$ presented below. Recalling that $B(Z \rightarrow \nu\nu) \simeq 18\%$, we may use Eq. (10b) to predict for every 17 events of the type $W^+ \rightarrow e^+ \nu$ produced directly, these should be about one monojet event arising from the process $M_c \rightarrow Z^0$ jet, $Z^0 \rightarrow \nu\nu$. This expectation is consistent with the UA1 data sample¹⁾. (We are using the term "monojet" somewhat loosely since our system X need not always decay into a single hadronic jet. In making comparisons with numbers of events, we have also ignored possible differences in acceptance and detection efficiencies for different topologies, e.g., $W \rightarrow e\nu$ vs. monojet plus missing transverse energy.)

Returning to Eq. (10b) and recalling that $B(Z^0 \rightarrow e^+e^-) \simeq 3\%$, we note that $M_c \rightarrow Z^0 X$ will also yield about one $Z^0 \rightarrow e^+e^-$ event for every 100 $W^+ \rightarrow e\nu$ events. This implies that about one $Z^0 \rightarrow e^+e^-$ event in the combined data samples of UA1 and UA2 may be due to $M_c \rightarrow Z^0 X$. The fraction of $Z^0 \rightarrow e^+e^-$ events arising from the mechanism $M_c \rightarrow Z^0 X$ is in the range 8 to 12%, not very different from the 8% fraction of $W^+ \rightarrow e^+ \nu$ events from $M_c \rightarrow WX$. Since the W 's and Z 's emerge from $M_c \rightarrow (W,Z)X$ with substantial momentum in the M_c rest frame (up to $|\vec{p}| \simeq 45$ GeV if $m_X = 0$), the bosons from this source would be produced at relatively large P_T and be accompanied by significant recoil hadron energy. As we mentioned earlier, there may be indications in the data for a second production mechanism for the intermediate vector bosons, with different properties from the Drell-Yan mechanism, more important for Z^0 than for W^+ . However, our mechanism, in its present form, does not imply a much higher fraction of "noisy" events in the Z^0 sample. This feature of the data remains a riddle, as does the frequency of monophoton events¹⁾.

The discussion above based on the spectator model shows that we are able to obtain rough consistency with the data on $e\nu$ jet, monojet, and "resonant" dijet events. To obtain hadronic branching fractions for M_c and to

relate these to $B(M_c \rightarrow WX)$ and to $B(M_c \rightarrow YX)$, we adopt a modification of techniques which have been generally useful for computations of exclusive processes⁸⁾ in QCD including partial widths of quarkonium systems. We take as an ansatz the separation of the decay amplitude into two terms. The first is the amplitude to have an off-shell $q\bar{q}$ state in the massive meson. The second is the hard scattering amplitude which transfers this $q\bar{q}$ state into the final particles on their mass-shells. This factorization is illustrated in Fig. 2a for $M_c \rightarrow q\bar{q}$, and in Fig. 2b for $M_c \rightarrow Vg$, where V denotes W , Z , g or γ . All exchange processes and mass effects associated with binding are included in the new physics "wave function" ψ . Both ψ and the pointlike hard scattering term T_H depend in principle upon a hadronic scale Q_0 at which the separation is made between the two types of interactions. The full amplitude is obtained as a convolution or loop-integral:

$$\mathcal{M}(p_i; Q_0) = \int \psi(q_i; Q_0) T_H(p_i, q_i; Q_0) \delta(q_i + q_k - p_k) \quad (11)$$

P_N is the four-vector momentum of M_c . In computing T_H we will neglect quark masses. Components in ψ containing more partons than $q\bar{q}$ are expected to yield corrections to Eq. (11) of order $\alpha_s(M_c^2)$ and of order $1/M_c^2$.

Because of the large momentum transfer involved, the process $M_c \rightarrow q\bar{q}$ is expected to proceed by single hard gluon exchange, as illustrated in Fig. 2a. Setting $q_i = z_i P$, we indeed find that the intermediate gluon carries the squared four-vector momentum $Q^2 \propto z_i M_c^2$, which is very large, justifying a perturbative approach^{1*)}. As in Ref. 8), we may introduce a distribution amplitude $\psi(z_i; Q_0)$ obtained by integrating ψ over the transverse components of q_i . Here the z_i are light-cone fractions: $q_i^+ = z_i P^+$. We then find

$$\mathcal{M}(M_c \rightarrow q\bar{q}) = \int \varphi(z_i) \delta(1 - z_1 - z_2) T_H(z_i; p_i) dz_1 dz_2 \quad (12)$$

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Large except, of course, when $z_i = 0$. However, ψ is expected to vanish as $z_i \rightarrow 0$ at least as fast as z_i , guaranteeing convergence of the integrals over z_i .

with

$$\Gamma_{\mu}^{q\bar{q}} = C_F \frac{g^2}{\sqrt{2}Q^2} \bar{u}(\phi_1) \delta_{\mu} (\not{p}_1 + \not{M}) \delta_5 \not{p}_2 U(\phi_2). \quad (13)$$

In Eq. (13), $C_F^{q\bar{q}}$ represents a colour matrix and the pseudoscalar nature of M_c is represented by the projection operator $(1/\sqrt{2})(\not{p} + \not{M})\gamma_5$. The normalization is such that the asymptotic form of $\phi(z_1)$ is $\phi(z_1, z_2) = f_{M_c} z_1 z_2$, where f_M is the "weak decay constant" of M_c (cf., for $\pi \rightarrow \mu\nu$, $f_{\pi} \approx 93$ MeV).

Carrying through the spin and colour sums, we derive

$$\Gamma(M_c \rightarrow q\bar{q}) = \frac{1}{2\pi M} g^4 \mathcal{I}_0. \quad (14)$$

with $\mathcal{I}_0 = \int dz \phi(z)/z(1-z)$, and $g^2 = 4\pi\alpha_s$. Adopting the asymptotic form mentioned above, we find $\mathcal{I}_0 = f_M$, and, with $\alpha_s = 0.1$, we can re-express Eq. (14) as

$$f_M \approx 2(M\Gamma(q\bar{q}))^{1/2}. \quad (15)$$

If $\Gamma(q\bar{q}) \approx 1$ GeV, we find $f_M \approx M/2\pi$, whereas $\Gamma(q\bar{q}) \approx 10$ GeV implies $f_M \approx M/2$. This range of values of $\Gamma(q\bar{q})$ brackets our earlier estimates and leads to a priori reasonable values of f_M for a hadronic state for which one expects f_M to be equal to M up to factors of order unity. The reasonable range of values obtained for f_M serves as a cross-check of our representation of M_c as an off-shell $q\bar{q}$ system.

Turning to the decay modes $M_c \rightarrow g\bar{g}$, $g\gamma$, gZ^0 and gW , we employ a similar perturbative approach now based on Fig. 2b. Owing to the symmetry of $\phi(z_1, z_2)$ under interchange of z_1 and z_2 , only the vector coupling terms survive for the W and Z^0 . We derive the amplitude

$$\Gamma_{\mu}^{gV} = C_F \frac{g g_V}{\sqrt{2}Q^2} \delta_i \epsilon_{\alpha\beta\mu\nu} p_1^{\alpha} p_2^{\beta} \epsilon_{\nu}^{\gamma} \epsilon_{\gamma}^{\delta} \quad (16)$$

where $Q^2 = -z(1-z)M_c^2 + 2M_V^2$, and g_V represents the appropriate vector coupling strength.

If M_c is an isovector system, $\Gamma(M_c \rightarrow g\bar{g}) = 0$, and we find

$$\Gamma(M_c \rightarrow g\gamma) = \frac{1}{\pi M} g^2 e^2 \mathcal{I}_0. \quad (17)$$

For an isoscalar state $M_c^{(0)}$, we obtain

$$\Gamma^{(0)}(M_c \rightarrow g\bar{g}) = \frac{3}{2\pi M} g^4 \mathcal{I}_0, \quad (18)$$

with $\Gamma^{(0)}(M_c \rightarrow g\gamma) = 1/9 \Gamma(M_c \rightarrow g\gamma)$. In Eq. (17), $e^2 = 4\pi\alpha_{em}$. The factor of three enhancement of $\Gamma(g\bar{g})$ relative to $\Gamma(q\bar{q})$ in Eqs. (14) and (18) is a combination of a colour factor of six reduced by a factor two associated with the identical boson final state.

In the cases of $M_c^0 \rightarrow gZ^0$ and $M_c^{\pm} \rightarrow gW^{\pm}$, we derive

$$\Gamma(M_c \rightarrow gW) = \frac{1}{\pi M} g^2 \left(\frac{e}{2\sqrt{2}S_W} \right)^2 \frac{8|\vec{p}_W|^3}{M^3} \mathcal{I}_W, \quad (19)$$

and

$$\Gamma(\nu_c \rightarrow g z^0) = \frac{1}{\pi M} g^2 \left(\frac{e}{4c_w s_w} \right)^2 \frac{8 |\vec{p}_z|^3}{M^2} \vec{p}_z^2. \quad (20)$$

In Eqs. (19) and (20), $s_w = \sin \theta_w$ and $c_w = \cos \theta_w$; $|\vec{p}|$ represents the magnitude of the three-vector momentum of the final vector boson in the M_c rest frame. We note that $2|\vec{p}_w|/M = (1-N_c^2/M^2) \approx 0.7$.

The integrals I_w and I_z which readily generalize I_0 in each case differ from I_0 because of non-trivial kinematic effects associated with the boson masses. The ratios $I_0/I_{w,z}$ cannot be calculated reliably in the framework of the collinear approximation followed here. Since we have dealt so far with highly off-shell quarks, we may consider the other extreme of the non-relativistic heavy quark limit for the processes gg, gZ, gW and $g\gamma$. In this case, one may easily obtain $I_0/I_{w,z} = 2|\vec{p}_{w,z}|/M$. We shall adopt this relationship in Eqs. (19) and (20) and in the Table, but consider the results obtained as valid at the order of magnitude level only.

For an isoscalar state, an additional factor of (1/9) would appear on the right-hand side of Eq. (20), and there are no gW decays.

Owing to contributions from gluon-gluon subprocesses, the production of an isoscalar state $M_c^{(0)}$ will be enhanced somewhat with respect to that for an isovector. At $M \approx 160$ GeV and $\sqrt{s} = 540$ GeV, this is not a large effect in $\bar{p}p$ collisions since, as mentioned before, the $\bar{u}u$ luminosity is considerably larger than the gg luminosity at this large mass. The strong suppression of the gZ^0 decay mode coupled with the large gg branching fraction [Eq. (18)] is, however, an undesirable aspect of an isoscalar state.

We summarize the above results for an isovector state [Eqs. (14), (17), (19) and (20)] in the second column of the Table (following the heavy quark approximation). We have set $\sin^2 \theta_w = 0.25$. With $\alpha_{em}/\alpha_s \approx 0.1$, the branching fractions in the Table imply that for each event of the type e^+e^- jet associated with $M_c \rightarrow W^+X$, there should be about one monojet event from $M_c \rightarrow Z^0X, Z^0 \rightarrow \nu\nu$; 20 γ -jet events from $M_c \rightarrow \gamma g$; and about 300 jet-jet events from the three charge states of M_c .

The rough equality of monojet and $e\nu$ jet rates, which does not depend on the heavy quark approximation, is consistent with the observations of the UA1 and UA2 collaborations. Most interesting perhaps is the prediction of a substantial number of γ -jet events, at a rate of the order γ -jet/jet-jet $\approx 1/15$, again independent of the approximations made for I_w and I_z . Since the probability for a typical QCD jet to emerge as a neutral electromagnetic shower is of the order of 10^{-3} , and comparable in rate to the prompt photon yield⁹⁾ at the values of p_T considered here, the γ -jet signal should be observed fairly readily above background. The data should be examined for this effect.

With the approximation made in connection with the ratio $I_0/I_{w,z}$, we would expect about 300 jet-jet events for each $e\nu$ -jet event. This is substantially larger than the signal of 69 ± 17 events discussed by the UA2 collaboration but can still be deemed acceptable. It is indeed not inconceivable that our crude models have led to an underestimate of the W and Z branching fractions.

When the energy of $\bar{p}p$ collisions is increased from $\sqrt{s} = 540$ GeV to $\sqrt{s} = 620$ GeV in the near future at CERN, and to $\sqrt{s} = 2$ TeV at Fermilab, interesting changes are expected in the ratio of the number of $e\nu$ jet and monojet events to $W^{\pm}$ and Z^0 events produced directly by the Drell-Yan mechanism. For a production mechanism based on $q\bar{q}$ subprocesses such as that explored in this letter, this ratio will grow by only $\sim 20\%$ from $\sqrt{s} = 540$ to 620 GeV and double from $\sqrt{s} = 540$ GeV to $\sqrt{s} = 2$ TeV. By contrast, if the unusual events at $M \approx 160$ GeV are generated principally from gluon-gluon subprocesses (as would be the case for Higgs, quarkonium, heavy quark or gluino production via $gg \rightarrow \tilde{g}\tilde{g}$) the ratio to direct $W^{\pm}$ events will grow by a factor of two from $\sqrt{s} = 540$ to 620 GeV, and by a factor of 70 from $\sqrt{s} = 0.54$ to 2 TeV. These dramatic differences will be easily distinguished.

Many of our qualitative conclusions are shared by models based on the production of excited quarks⁵⁾. It is probably impossible at present to test differences such as possible charge correlations associated with the fact that a gluon recoil in our case replaces the quark recoil of the excited quark model. There is some difference in the predicted energy dependences, but it

is not significant over the range $\sqrt{s} = 540$ to 620 GeV. One salient difference is that if octet states are present at $M \sim 160$ GeV, there is no reason not to also expect $\bar{3}$ diquark systems which would be in the vicinity of 70 GeV, following the mass ratios of the crude model referred to above. These would lead to a related anomaly in the jet-jet decay channel.

Massive colour octet mesons ($M \sim 250$ GeV) have also been predicted in versions of extended technicolour¹⁰⁾. In this letter we have refrained from committing ourselves to new dynamical substructure, but much of our phenomenological analysis would apply directly to the production and decay of technicolour states. The calculational procedures we adopted would be readily justified because of the large mass of the relevant technifermions. A possible difficulty for technicolour, however, is that a plethora of other states is also predicted, including colour singlet technipions with masses lower than those of the colour octet states. Any difficulties with an isoscalar state excited through gg subprocesses would also apply.

In conclusion, we reiterate that the speculative nature of our considerations cannot be minimized. Nevertheless, there is limited freedom in the estimates of rates and branching fractions, and the general agreement with data is noteworthy. Increased statistics in the forthcoming collider runs will be telling. A thorough study of the γ -jet channel in the present data sample may be very rewarding.

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TABLE
 Branching ratios for an isovector state M_c , relative to that for $M_c \rightarrow q\bar{q}$.

Decay Modes	Branching Ratios	Observation
$q\bar{q}$	1	Peak in the jet-jet mass distribution
$q\gamma$	$2\alpha_{em}/\alpha_s$	Hadron jet balanced by a single photon
gW	$2 \vec{p}_W (\frac{\alpha_{em}}{\alpha_s})(\frac{1}{M})$	Jet (or jets) plus e (or μ) and missing energy from $W \rightarrow \ell\nu$
gZ^0	$2 \vec{p}_Z (\frac{2\alpha_{em}}{3\alpha_s})(\frac{1}{M})$	Monojet plus missing energy from $Z^0 \rightarrow \nu\nu$, 18% of the time, or monojet plus $\ell^+\ell^-$ from $Z^0 \rightarrow \ell^+\ell^-$

FIGURE CAPTIONS

Fig. 1: Simple spectator model diagram for the decay $M_c \rightarrow WX$.

Fig. 2: a) Diagram illustrating the decay of M_c into a quark-antiquark final state via single gluon exchange. Thick lines represent off-shell quarks.

b) Diagram illustrating the decay of M_c into a gluon, g , with momentum p_1 and a vector/axial vector with momentum p_2 (dashed line represents Z , W , g or γ); g and g_v are coupling strengths.

Fig. 1

Fig. 2 (a)

Fig. 2 (b)