

SEARCH FOR NEW QUARKS SUGGESTED BY THE SUPERSTRING

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We discuss the possible phenomenology of the additional charge $(-\frac{1}{3})$ colour triplet particles $D_{1/2}$ of spin $\frac{1}{2}$ and D_0, \bar{D}_0^c of spin 0 which are contained in each matter generation if the superstring is compactified on a Calabi-Yau manifold. In the minimal model with no intermediate mass scales and supersymmetry breaking fed into the observable sector via a gaugino mass, either a spin- $\frac{1}{2}$ or a spin-0 D particle could be the lightest, and either could be as light as the present lower bound from e^+e^- experiments of about 20 GeV. The D particles could behave as leptoquarks coupling to quarks and leptons, in which case the single production process $ep \rightarrow D_0(\bar{D}_0^c) + X$ would occur, and $D_0(\bar{D}_0^c) \rightarrow \ell q, D_{1/2} \rightarrow \ell q \tilde{\chi}$ decays would dominate, where $\tilde{\chi}$ is the lightest supersymmetric particle. Alternatively, the D particles could behave as diquarks coupling to pairs of antiquarks, in which case the single production process $(\bar{p})p \rightarrow D_0(\bar{D}_0^c, \bar{D}_0, D_0^c) + X$ would occur, and $D_0(\bar{D}_0^c) \rightarrow \bar{q}\bar{q}, D_0^c(\bar{D}_0) \rightarrow qq, D_{1/2} \rightarrow \bar{q}\bar{q}\tilde{\chi}$ decays would dominate. We present cross sections for $ep \rightarrow D_0(\bar{D}_0^c) + X, (\bar{p})p \rightarrow D_0(\bar{D}_0^c, D_0^c, \bar{D}_0) + X, e^+e^- \rightarrow D_0\bar{D}_0(D_0^c\bar{D}_0^c)$ and $D_{1/2}\bar{D}_{1/2}, (\bar{p})p \rightarrow D_0\bar{D}_0(D_0^c\bar{D}_0^c)$ and $D_{1/2}\bar{D}_{1/2} + X$. We calculate the experimental signals and estimate backgrounds for D production and decay in these processes. The decays $D_0(\bar{D}_0^c) \rightarrow \ell q$ and $D_{1/2} \rightarrow \ell q \tilde{\chi}$ or $\bar{q}\bar{q}\tilde{\chi}$ would be detectable in most of these reactions, but $D_0(\bar{D}_0^c) \rightarrow \bar{q}\bar{q}$ decays may only be detectable in $e^+e^- \rightarrow D_0\bar{D}_0(D_0^c\bar{D}_0^c)$ collisions.

1. Introduction

The low-energy phenomenology of the superstring [1] depends on the compactification scheme which is adopted. An attractive proposal which avoids many theoretical and phenomenological pitfalls is that the six unseen dimensions are compactified on a Kähler manifold of SU(3) holonomy and zero torsion [2]. In this case the observable particles appear in generations containing 27 fields with the quantum numbers of the fundamental representation of E_6 , plus possible split multiplets coming from $\underline{27} + \overline{\underline{27}}$ representations of E_6 , and the four-dimensional observable gauge group is some rank 5 or 6 subgroup of E_6 [3]. In some models this gauge

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group is broken by the Higgs mechanism, and the light particles do not include all the fields in a $\underline{27}$ of E_6 . However, it is quite possible that some exotic matter particles survive down to low energies. Thus one can expect in many of the conceivable symmetry-breaking patterns to see additional matter particles and/or gauge interactions at low energies. Although they have not yet been worked out in detail, some orbifold compactifications [4] could offer similar phenomenological possibilities [5].

Constraints on additional neutral current effects at low energies have been discussed extensively [6, 7], as well as possible signals of new gauge bosons at future colliders [8]. The possible phenomenology of colour-singlet superstring particles has also been explored [9]. This paper is concerned with the possible signatures of additional colour-triplet superstring particles. These have also been previously discussed [10], though the emphasis of previous studies has been somewhat different from that taken in this work.

We start with a pedagogical discussion of the particle content and of the possible interactions of the 27 fields in the fundamental representation of E_6 . The $\underline{27}$ representation splits up under $SO(10)$ to $\underline{16} + \underline{10} + \underline{1}$ and then under $SU(5)$ to $(\underline{10} + \underline{\bar{5}} + \underline{1}) + (\underline{5} + \underline{\bar{5}}) + \underline{1}$. According to the conventional assignment, the usual quarks and leptons are members of the first $\underline{10}$ and $\underline{\bar{5}}$. The states of the additional $\underline{5}$ and $\underline{\bar{5}}$ can be identified with the usual Higgs doublets H and \bar{H} , and as new additional colour triplets D and D^c . There is also an additional $SO(10)$ and $SU(5)$ singlet called N , and an $SU(5)$ (but not $SO(10)$) singlet called ν^c .

The decomposition of the $\underline{27}$ representation under the standard model subgroup $G_0 \equiv SU(3)_C \times SU(2)_L \times U(1)_Y$ consists of the following left-handed multiplets:

$$Q \equiv \begin{pmatrix} u \\ d \end{pmatrix} \sim (3, 2, +\frac{1}{6}), \quad u^c \sim (\bar{3}, 1, -\frac{2}{3}), \quad e^c \sim (1, 1, +1),$$

$$D \sim (3, 1, -\frac{1}{3}), \quad H \equiv \begin{pmatrix} H^+ \\ H^0 \end{pmatrix} \sim (1, 2, +\frac{1}{2}); \quad (1.1)$$

$$d^c, D^c \sim (\bar{3}, 1, +\frac{1}{3}), \quad L \equiv \begin{pmatrix} \nu \\ e \end{pmatrix}, \quad \bar{H} \equiv \begin{pmatrix} \bar{H}^0 \\ \bar{H}^- \end{pmatrix} \sim (1, 2, -\frac{1}{2}),$$

$$N, \nu^c \sim (1, 1, 0). \quad (1.2)$$

The first two numbers in brackets denote the dimensions of the $SU(3)_C$ and $SU(2)_L$ representations respectively, and the third one the $U(1)_Y$ hypercharge. The above states have, besides the usual gauge couplings, also generalized Yukawa interactions. The most general superpotential couplings allowed by the Hosotani [11] symmetry-breaking mechanism, which is the only one presently available for breaking E_6

down to an acceptable subgroup at the compactification scale, are [3]:

$$f = h_U Q u^c H + h_E L e^c \bar{H} \quad (1.3a)$$

$$+ h_\nu H L \nu^c + \lambda H \bar{H} N \quad (1.3b)$$

$$+ \lambda_\nu D d^c \nu^c + k D D^c N \quad (1.3c)$$

$$+ \lambda_L Q D^c L + h_D Q d^c \bar{H} \quad (1.3d)$$

$$+ \lambda_c D e^c u^c + \lambda_Q D Q Q + \lambda_c D^c u^c d^c. \quad (1.3e)$$

In a conventional GUT, the couplings $h_U, h_E, h_\nu, \lambda \dots$ would be related by Clebsch-Gordan coefficients of order unity, and therefore the new states D and D^c would need to be very heavy in order to avoid rapid proton decay [12]. The new feature of the string compactification scenario is that not only is the GUT gauge group broken by the Hosotani [11] mechanism, but also the Clebsch-Gordan relations between Yukawa couplings are broken [3]. This leads to the phenomenologically interesting possibility that the new colour triplets D and D^c could be light and behave either as leptoquarks or as diquarks.

The D and D^c particles acquire their mass through the vacuum expectation value of the N field. The fermionic members of the D/D^c supermultiplets will combine into Dirac fermions of mass $m_D = k \langle 0|N|0 \rangle$, denoted here by $D_{1/2}$. The masses of their scalar partners D_0, D_0^c receive additional contributions from soft super-symmetry-breaking terms which are of the same order as conventional Higgs or squark masses, and hence are expected to be $O(m_w$ to 1 TeV). It is possible that $\langle 0|N|0 \rangle$ is very large, but in that case the $H\bar{H}N$ coupling (1.3b) would need to be very small in order for $\langle 0|H, \bar{H}|0 \rangle$ to be $O(m_w)$ as required. It is therefore interesting to study the case that $\langle 0|N|0 \rangle$ is not very much larger than m_w , where one expects both the $D_{1/2}$ fermions and the D_0, D_0^c bosons to weigh $\leq O(1)$ TeV and hence be accessible to experiment. *A priori*, one could have mixing between d^c and D^c and between L and \bar{H} . We discuss this problem carefully in sect. 2 and argue that in the context of a minimal model there are no rotations between these sectors to be considered (for an alternative point of view, see ref. [10]).

The non-observation of supersymmetric particles and the absence of additional neutral current effects constrain the parameters of superstring models and hence the possible masses of the $D_{1/2}$ and D_0, D_0^c particles. We analyze these constraints in sect. 2, working in the context of a minimal rank-5 low-energy model [6], where the dominant source of supersymmetry breaking in the observable sector is a gaugino mass term. We also incorporate more model-dependent constraints which follow from dynamical calculations of radiative electroweak breaking in this special theoretical framework. We find that some of the $D_{1/2}$ and D_0, D_0^c particles could be

as light as the present bounds on new particles in e^+e^- collisions: $m \geq 20$ GeV. The D particles could be produced by single or pair production mechanisms. Single production could proceed through quark–quark annihilation in hadron–hadron collisions or quark–lepton annihilation in ep collisions through the diquark and leptoquark couplings of eq. (1.3). The two processes cannot happen simultaneously, because otherwise rapid proton decay could be mediated by D_0 exchange. The couplings of the D particles which influence their possible production mechanisms and decay signatures are discussed in sect. 3. The cross sections for these two single-production mechanisms are calculated and compared with conventional backgrounds. We find that the process $ep \rightarrow (D_0/\bar{D}_0^c \rightarrow q\ell) + X$ could be observable at HERA if the D_0 mass is below ~ 250 GeV. The diquark annihilation process which could take place at the hadron–hadron colliders, $(\bar{p})p \rightarrow D_0/D_0^c \rightarrow \bar{q}\bar{q}/qq + X$ would, however, be overwhelmed by the conventional QCD jet–jet background. However, the pair-production process could be a source of observable D particles at hadron–hadron colliders. In sect. 4 we calculate cross sections for the possible pair-production mechanisms $hh \rightarrow D_0\bar{D}_0(D_0^c\bar{D}_0^c) + X$ and $D_{1/2}\bar{D}_{1/2} + X$ which proceed through the conventional strong interactions, and explore the possible event signatures. Since the gauge couplings of D particles are the same as those of usual quarks, these cross-sections are comparable to those for conventional squarks or t quarks respectively. The decays $D_0/\bar{D}_0^c \rightarrow q + \nu$, $\bar{D}_0/D_0^c \rightarrow \bar{q} + \nu$, $D_{1/2} \rightarrow q + \nu + \tilde{\chi}$ and $D_{1/2} \rightarrow \bar{q} + \bar{q} + \tilde{\chi}$, where $\tilde{\chi}$ denotes the lightest supersymmetric particle, have missing energy signatures analogous to those expected for conventional squarks and gluinos [13]. The present UA1 missing energy search [14] may be able to exclude D_0 or $D_{1/2}$ weighing up to $O(60)$ GeV, but this needs to be investigated by the experimentalists themselves. Any such limit could be improved by future experiments using ACOL at CERN, or the Tevatron Collider. The decays $D_{1/2} \rightarrow q + \ell + \tilde{\chi}$ give experimental signatures with leptons, jets and missing energy which could have shown up in the UA1 dimuon sample [15] or in their top search, and again present data might have sensitivity for $m_{D_{1/2}} \lesssim 60$ GeV. The decays $D_0(\bar{D}_0^c) \rightarrow \bar{q} + \bar{q}$ and $\bar{D}_0(D_0^c) \rightarrow q + q$ would give four-jet final states which would be drowned by the QCD four-jet background, but the decays $D_0/\bar{D}_0^c \rightarrow q\ell$ ($\bar{D}_0/D_0^c \rightarrow \bar{q}\ell$) would give dilepton + dijet signatures which should be easier to detect, and could show up in the UA1 dimuon sample [15]. A useful limit on them could perhaps be obtained with present data.

Finally, sect. 5 presents some conclusions and proposes some possible directions for future experiments which could search for D particles.

2. Constraints on the (D_0, D_0^c) and $D_{1/2}$ masses

Here we discuss in detail the mass matrices for the exotic quarks $D_{1/2}$ and squarks D_0, D_0^c in the realistic case of $n_G = 3$ generations. Where it is necessary to be specific, we will work in the framework of the minimal rank-5 gauge group [6]

which can be provided by Hosotani symmetry breaking. However, the following analysis could be extended to any similar model with light exotic quarks and squarks.

First we discuss the possibility mentioned in sect. 1 of mixings between the fields d^c and D^c , L and \bar{H} , ν^c and N , which are seen from eq. (1.2) to have identical $SU(3)_C \times SU(2)_L \times U(1)_Y$ quantum numbers. If the low-energy group commutes with $SU(2)_N$, which is the case for the minimal model, we may without loss of generality identify L with the standard lepton doublet, in which case lepton number conservation requires $\langle 0|\tilde{\nu}|0\rangle = 0$. In order to identify the mass eigenstates corresponding to the ordinary charge $-\frac{1}{3}$ quarks and charged leptons, one must examine the corresponding mass matrices, which take the forms

$$(d, D)M_{1/3} \begin{pmatrix} d^c \\ D^c \end{pmatrix}: \quad M_{1/3} = \begin{pmatrix} h_D \langle 0|\bar{H}^0|0\rangle & 0 \\ \lambda_\nu \langle 0|\tilde{\nu}^c|0\rangle & k \langle 0|N|0\rangle \end{pmatrix}, \quad (2.1)$$

$$(e^c, H^+)M_1 \begin{pmatrix} e \\ \bar{H}^- \end{pmatrix}: \quad M_1 = \begin{pmatrix} h_E \langle 0|\bar{H}^0|0\rangle & 0 \\ h_\nu \langle 0|\tilde{\nu}^c|0\rangle & \lambda \langle 0|N|0\rangle \end{pmatrix}. \quad (2.2)$$

The physical squared masses are given by the eigenvalues of the matrices $M_{1/3}^+ M_{1/3}$ and $M_1^+ M_1$ respectively. We assume for simplicity real Yukawa coupling constants and vacuum expectation values (VEVs). Note that, in order to avoid unacceptable zero mass eigenvalues in (2.1) and (2.2), it must be that $\langle 0|N|0\rangle \neq 0$. On the other hand, the possibilities that $\lambda^D \langle 0|\tilde{\nu}^c|0\rangle, \lambda^L \langle 0|\tilde{\nu}^c|0\rangle \neq 0$ have potential phenomenological problems with flavour-changing neutral currents and the physical values of m_d, m_e [17]. Thus we will assume in the rest of this paper that $\langle 0|\tilde{\nu}^c|0\rangle = 0$, which implies $\langle 0|\bar{H}^0|0\rangle \neq 0$ to avoid zero eigenvalues in (2.1) and (2.2). It must also be that $\langle 0|H^0|0\rangle \neq 0$ to give masses to the quarks of charge $\pm \frac{2}{3}$. We are then led to a situation where we can identify (d, d^c) with the ordinary quarks (of mass $h_D \langle 0|\bar{H}^0|0\rangle$) and (D, D^c) with the new exotic quarks (of mass $k \langle 0|N|0\rangle$), and there is no mixing between the two sectors. Similarly, we will identify (e, e^c) with the ordinary leptons, and (H^+, \bar{H}^-) with the charged Higgs fields.

In the case of a three-generation model, the fields which can acquire non-vanishing VEVs are H_a^0, \bar{H}_a^0 and N_a , where $a = 1, 2, 3$ is a generation index. However, it is always possible [18] to redefine those fields, via independent rotations in generation space, in such a way that $\langle 0|H_a^0|0\rangle = \langle 0|\bar{H}_a^0|0\rangle = \langle 0|N_a|0\rangle = 0$ for $a = 1, 2$ and the only non-vanishing VEVs are:

$$\langle 0|H_3^0|0\rangle \equiv v, \quad \langle 0|\bar{H}_3^0|0\rangle \equiv \bar{v}, \quad \langle 0|N|0\rangle \equiv x. \quad (2.3)$$

In this new basis, the part of the superpotential contributing to the $D_{1/2}$ masses is

simply:

$$f_{D_{1/2}}^{\text{mass}} = k_{ab3} D_a D_b^c N_3. \quad (2.4)$$

We have not used, up to now, the freedom of redefining the fields D_a and D_a^c by independent rotations in generation space. We can therefore move to a convenient basis where the following relations hold:

$$k_{ab3} = 0 \quad \text{for } a \neq b. \quad (2.5)$$

Assuming, without loss of generality, that $x > 0$, the two-component spinors of each generation contained in the chiral superfields D_a, D_a^c will then combine to form three Dirac spinors $D_{1/2a}$ of masses

$$m_{D_a} = k_{aa3} x \quad (a = 1, 2, 3). \quad (2.6)$$

The mass matrix involving the scalar components of D_a and D_a^c has a slightly more complicated structure. Even in the convenient basis we chose above, one has in general a real symmetric 6×6 matrix, with non-trivial intergenerational mixing. The matrix elements are

$$\begin{aligned} M_{D_1 \bar{D}_1}^2 &= \tilde{m}_{D_{11}}^2 + m_{D_1}^2 + G^2(D), & M_{\bar{D}_1^c D_1^c}^2 &= \tilde{m}_{D_{11}^c}^2 + m_{D_1}^2 + G^2(D^c), \\ M_{D_2 \bar{D}_2}^2 &= \tilde{m}_{D_{22}}^2 + m_{D_2}^2 + G^2(D), & M_{\bar{D}_2^c D_2^c}^2 &= \tilde{m}_{D_{22}^c}^2 + m_{D_2}^2 + G^2(D^c), \\ M_{D_3 \bar{D}_3}^2 &= \tilde{m}_{D_{33}}^2 + m_{D_3}^2 + G^2(D), & M_{\bar{D}_3^c D_3^c}^2 &= \tilde{m}_{D_{33}^c}^2 + m_{D_3}^2 + G^2(D^c), \\ M_{D_1 D_1^c}^2 &= \eta_{113}^k x + \sum_c \lambda_{33c} k_{11c} \nu \bar{\nu}, & M_{D_1 \bar{D}_2}^2 &= \tilde{m}_{D_{12}}^2, \\ M_{D_1 D_2^c}^2 &= \eta_{123}^k x + \sum_c \lambda_{33c} k_{12c} \nu \bar{\nu}, & M_{D_1 \bar{D}_3}^2 &= \tilde{m}_{D_{13}}^2, \\ M_{D_1 D_3^c}^2 &= \eta_{133}^k x + \sum_c \lambda_{33c} k_{13c} \nu \bar{\nu}, \\ M_{\bar{D}_1^c \bar{D}_2}^2 &= \eta_{213}^k x + \sum_c \lambda_{33c} k_{21c} \nu \bar{\nu}, & M_{\bar{D}_1^c D_2^c}^2 &= \tilde{m}_{D_{21}^c}^2, \\ M_{\bar{D}_1^c \bar{D}_3}^2 &= \eta_{313}^k x + \sum_c \lambda_{33c} k_{31c} \nu \bar{\nu}, & M_{\bar{D}_1^c D_3^c}^2 &= \tilde{m}_{D_{31}^c}^2, \\ M_{D_2 D_2^c}^2 &= \eta_{223}^k x + \sum_c \lambda_{33c} k_{22c} \nu \bar{\nu}, & M_{D_2 \bar{D}_3}^2 &= \tilde{m}_{D_{23}}^2, \\ M_{D_2 D_3^c}^2 &= \eta_{233}^k x + \sum_c \lambda_{33c} k_{23c} \nu \bar{\nu}, \\ M_{\bar{D}_2^c \bar{D}_3}^2 &= \eta_{323}^k x + \sum_c \lambda_{33c} k_{32c} \nu \bar{\nu}, & M_{\bar{D}_2^c D_3^c}^2 &= \tilde{m}_{D_{32}^c}^2, \\ M_{D_3 D_3^c}^2 &= \eta_{333}^k x + \sum_c \lambda_{33c} k_{33c} \nu \bar{\nu}. \end{aligned} \quad (2.7)$$

Symbols appearing in the above formula correspond to the following general expressions for the superpotential and the soft supersymmetry-breaking part of the scalar potential:

$$f = \lambda_{abc} H_a \bar{H}_b N_c + k_{abc} D_a D_b^c N_c + \dots, \quad (2.8)$$

$$V_{\text{soft}} = \tilde{m}_{D_a}^2 \bar{D}_a^* D_b + \tilde{m}_{D_b}^2 \bar{D}_a^c D_b^c + \dots + \eta_{abc}^k (D_a D_b^c N_c + \text{h.c.}) + \dots, \quad (2.9)$$

where the dots represent terms which do not contribute to (2.7), m_{D_a} are the fermion masses given by eq. (2.6) and $a, b, c = 1, 2, 3$. Moreover, $G^2(D)$ and $G^2(D^c)$ are the generation-independent D -term contributions to the scalar masses:

$$G^2(D) = \frac{5}{9} \frac{m_W^2 \tan^2 \theta_W}{1 + (\bar{v}/v)^2} \left[1 + \left(\frac{\bar{v}}{v} \right)^2 - 2 \left(\frac{x}{v} \right)^2 \right], \quad (2.10)$$

$$G^2(D^c) = \frac{5}{9} \frac{m_W^2 \tan^2 \theta_W}{1 + (\bar{v}/v)^2} \left[1 - \frac{1}{2} \left(\frac{\bar{v}}{v} \right)^2 - \frac{1}{2} \left(\frac{x}{v} \right)^2 \right]. \quad (2.11)$$

To bring the discussion to a reasonable degree of simplicity, some short-cuts are necessary. Our first simplification is to assume that the only Yukawa couplings contributing significantly to the renormalization group equations (RGE) for the soft supersymmetry-breaking parameters are those diagonal in generation space ($\lambda_{aaa}, k_{aaa}, \dots$). In this case, the soft supersymmetry-breaking part of the scalar potential assumes the simplified form:

$$V_{\text{soft}} = \tilde{m}_{D_a}^2 |D_a|^2 + \tilde{m}_{D_c}^2 |D_a^c|^2 + \dots \\ + k_{abc} A_{abc}^k (D_a D_b^c N_c + \text{h.c.}) + \dots. \quad (2.12)$$

Moreover, the mass matrix (2.7) for D_0 squarks becomes block-diagonal in generation space, with each 2×2 block given by

$$\mathcal{M}_{0_a}^2 = \begin{pmatrix} \tilde{m}_{D_a}^2 + m_{D_a}^2 + G^2(D) & m_{D_a} A_{aa3}^k + \lambda_{333} k_{aa3} v \bar{v} \\ m_{D_a} A_{aa3}^k + \lambda_{333} k_{aa3} v \bar{v} & \tilde{m}_{D_a}^2 + m_{D_a}^2 + G^2(D^c) \end{pmatrix}. \quad (2.13)$$

Let us examine now the phenomenological constraints on the quantities appearing in eqs. (2.6) and (2.13). These constraints are derived under the plausible assumption that the main source of supersymmetry breaking in the observable sector is a universal gaugino mass $m_{1/2}$ [19], assigned at the grand unification scale M_X , which in turn generates all the other SUSY-breaking parameters via radiative corrections, taken into account by the renormalization group equations.

TABLE 1
Typical values [6] of the model parameters, specified at the electroweak symmetry-breaking scale

0.2	\lesssim	\bar{v}/v	\lesssim	0.6
2.8	\lesssim	x/v	\lesssim	10
100 GeV	\lesssim	$m_{1/2}$	\lesssim	500 GeV
0.15	\lesssim	λ_{333}	\lesssim	0.35
3.00	\lesssim	$\tilde{m}_{D_3}^2/m_{1/2}^2$	\lesssim	3.35
2.95	\lesssim	$\tilde{m}_{D_3^c}^2/m_{1/2}^2$	\lesssim	3.30
-3.5	\lesssim	$A_{333}^k/m_{1/2}$	\lesssim	-3.0
0.25	\lesssim	k_{333}	\lesssim	0.55

The couplings k_{aa3} ($a = 1, 2$) are assumed to be significantly smaller than the coupling k_{333} .

Of the three VEVs v , \bar{v} and x , only two combinations are independent since the constraint

$$m_w = \frac{e}{\sin \theta_w} \sqrt{\frac{v^2 + \bar{v}^2}{2}} \approx 82 \text{ GeV}$$

must be satisfied, where e is the running electric charge evaluated at $Q^2 \sim m_w^2$: it is then convenient to have as independent parameters x/v and \bar{v}/v .

As other parameters appearing in (2.6) and (2.13), it is convenient to take $m_{1/2}$, $(\tilde{m}_{D_a}/m_{1/2})^2$, $(\tilde{m}_{D_a^c}/m_{1/2})^2$, $A_{aa3}^k/m_{1/2}$, λ_{333} and k_{aa3} . To simplify the discussion further, we shall assume that the only non-negligible Yukawa couplings contributing to the RGEs are indeed $\lambda_{333} \equiv \lambda$, $k_{333} \equiv k$ and $h_{U_{333}} \equiv h$ (where h provides the top quark mass), so that the RGEs can be written down in the simplified form used in ref. [6]. In this case, indicative ranges of variation for the different parameters are presented in table 1. The lower limits on x/v and $m_{1/2}$ are derived from the present limits on sparticle masses and on the mass of a second neutral gauge boson. The corresponding upper limits resemble a form of naturalness constraint, and are more subjective. The ranges of variation of \bar{v}/v , λ_{333} and k_{333} are obtained from a detailed dynamical study of the effective scalar potential at low energy, solving numerically the RGEs for the soft SUSY-breaking parameters. The ranges of variation of $(\tilde{m}_{D_3}/m_{1/2})^2$, $(\tilde{m}_{D_3^c}/m_{1/2})^2$ and $A_{333}^k/m_{1/2}$ correspond to the solutions of the RGEs for allowed values of λ_{333} , k_{333} and of the top mass. The upper limits (in absolute value) correspond to small values of k_{333} , and vice versa*. The Yukawa coupling k_{333} must be greater than ~ 0.25 to generate the desired hierarchy $x/v \gtrsim 2.8$, and cannot be larger than ~ 0.55 if one wants to avoid charge-breaking minima. The other couplings k_{113} and k_{223} are assumed to be smaller than k_{333} , and there are no lower limits on them apart from the experimental limits on m_{D_a} .

* For $a = 1, 2$ we expect the corresponding quantities to be close to these limits, according to our assumption that $k_{113, 223} \ll k_{333}$.

Let us see now what are the indications on the $D_{1/2}$ and D_0 masses coming from explicit model calculations. To begin with, let us consider the $D_{1/2}$ masses, given by eq. (2.6), which in terms of the parameters of table 1 reads:

$$m_{D_a} = k_{aa3}x = k_{aa3}\left(\frac{x}{v}\right)\left\{2\sqrt{2}G_F\left[1 + (\bar{v}/v)^2\right]\right\}^{-1/2}, \quad (2.14)$$

where $G_F = 1.166 \times 10^{-5} \text{ GeV}^{-2}$ is the Fermi constant. Using the limits of table 1, one finds $m_{D_3} \geq 105 \text{ GeV}$, while for $a = 1, 2$ the lower limit on m_{D_a} is the experimental one:

$$m_{D_a} \geq 20 \text{ GeV} \quad (a = 1, 2). \quad (2.15)$$

An indicative upper limit on the $D_{1/2}$ masses comes from the naturalness constraint, which forbids very large values of x/v , since they would require a fine tuning of the parameters in the scalar potential which would be unstable under radiative corrections [6]. Using the values of table 1, one finds:

$$m_{D_1}, m_{D_2} < m_{D_3} \leq 1 \text{ TeV}. \quad (2.16)$$

In summary, model calculations suggest for the $D_{1/2}$ quarks masses ranging from 20 GeV to 1 TeV. Although the scale of $D_{1/2}$ masses is generally the same as the scale of the extra Z' boson, since both are related to the breaking of the extra $U(1)_E$ gauge factor, some $D_{1/2}$ particles may be considerably lighter than the Z' if they have small Yukawa couplings to the N field which acquires a VEV.

Let us comment now on the structure of the D_0 mass matrices, eq. (2.13). First of all, note that, for the range of parameters given in table 1, always $m_{D_a}A_{aa3}^k \gg \lambda_{333}k_{aa3}v\bar{v}$, so that we can neglect this last contribution to the off-diagonal entry. Note also that, for large values of x/v , the D -term contributions of eqs. (2.10) and (2.11) become negative and big. On the other hand, the positive contributions $\tilde{m}_{D_a}^2$ and $\tilde{m}_{D_a^c}^2$ are proportional to $m_{1/2}^2$. This suggests that light D_0 masses can be obtained for relatively high values of (x/v) and relatively small values of $m_{1/2}$. However, detailed model calculations [6] exhibit a strong correlation between x/v and $m_{1/2}$; in particular, to avoid $(\text{mass})^2$ for the charged “unhiggses” (H_a^+ , \bar{H}_a^-), ($a = 1, 2$), which are too small or negative, a relation of the type $m_{1/2} > K_{\text{cr}}(x/v)$ must be satisfied, with K_{cr} slightly dependent on the other parameters, and generically in the range $K_{\text{cr}} \approx 55 \div 65 \text{ GeV}$. Even taking into account this last constraint, together with the correlation between $\tilde{m}_{D_a}^2$, $\tilde{m}_{D_a^c}^2$ and A_{aa3}^k , one finds the following results:

(i) Calling D_{0a}^1 and D_{0a}^2 the two eigenvalues of the mass matrix (2.13), with masses $m_{D_{0a}^1} \leq m_{D_{0a}^2}$, the following two hierarchies are both possible:

$$m_{D_a} < m_{D_{0a}^1} < m_{D_{0a}^2}, \quad (2.17a)$$

$$m_{D_{0a}^1} < m_{D_a} < m_{D_{0a}^2}. \quad (2.17b)$$

(ii) Under our set of assumptions, masses of the D_{03}^1 ($D_{01,2}^1$) squarks as low as 100 GeV (170 GeV) are allowed. These lower limits come from the study of the “unhiggs” mass spectrum, and have been derived under some model-dependent assumptions [6]. Relaxing these assumptions, they can be significantly lowered and could be as light as 20 GeV. On the other hand, the addition of a common scalar mass to the primordial soft SUSY-breaking terms would generally lead to higher D_{0a} masses, favouring the case of eq. (2.17a) with respect to the one of eq. (2.17b).

(iii) When the lowest possible values for the D_{0a}^1 masses are attained, the mass eigenstate D_{0a}^1 is always an admixture, with contributions of almost equal magnitude, of the interaction eigenstates D_a and \bar{D}_a^c .

This is essentially all the information on $D_{1/2}, D_0$ masses that one is able to extract from model calculations. Despite all the constraints on the different parameters, there is still a large freedom for the resulting spectrum of $D_{1/2}, D_0$ masses, which is not significantly bounded from below, thus motivating the phenomenological analysis of the following sections.

3. Single D production and decay

3.1. D COUPLINGS

As was mentioned in the introduction, the D, D^c supermultiplets may couple directly to ordinary quarks and leptons through the following superpotential terms:

$$f \ni \lambda_Q D Q Q + \lambda_c D^c u^c d^c \quad (3.1a)$$

$$+ \lambda_L D^c L Q + \lambda_e D u^c e^c \quad (3.1b)$$

$$+ \lambda_\nu D d^c \nu^c. \quad (3.1c)$$

No other couplings are allowed by invariance under the E_6 Cartan subalgebra, which is maintained by the Hosotani gauge symmetry-breaking mechanism. In order to avoid rapid proton decay, the first row of couplings (3.1a) cannot be present simultaneously with the second (3.1b) and third (3.1c) rows of possible terms in f . Topological zeroes and/or discrete symmetries could in principle forbid the disallowed combinations of couplings, though no explicit example has yet been exhibited. Dirac neutrino masses can be made to vanish naturally only if the couplings (λ_L, λ_e) of (3.1b) are not present simultaneously with the coupling λ_ν of (3.1c), and the same remarks about topological zeroes and/or discrete symmetries apply. We will take the sets (3.1a), (3.1b) and (3.1c) as alternative cases to study.

3.2. COUPLINGS TO QUARKS

As discussed earlier, in the class of models we consider, no mixing of the D particles with conventional charge $-\frac{1}{3}$ quarks is possible. Single production of the

scalar D is possible in hadron-hadron collisions via $\bar{q} + \bar{q} \rightarrow D_0$ ($q + q \rightarrow D_0^c$), followed by $D_0 \rightarrow \bar{q} + \bar{q}$ ($D_0^c \rightarrow q + q$). The squared amplitudes $|\mathcal{M}|^2$ for the decays of D_0 and D_0^c to quark pairs are given by

$$|\mathcal{M}_{D_0}|^2 = 24\lambda_Q^2 m_{D_0}^2, \quad |\mathcal{M}_{D_0^c}|^2 = 6\lambda_c^2 m_{D_0^c}^2, \quad (3.2)$$

and we have no idea what the ratio of the couplings λ_Q and λ_c might be. The differential cross section for the process $h + h \rightarrow (D_0 \rightarrow qq) + X$ is given by

$$d^2\sigma = (1/48\pi)\hat{s} \{ u(x_1)d(x_2) + u(x_2)d(x_1) + \dots \} dx_1 dx_2 \delta(sx_1x_2 - \hat{s}) \\ \times \left\{ \frac{\lambda_Q^4(64/9)}{\left((\hat{s} - m_{D_0}^2)^2 + \Gamma_{D_0}^2 m_{D_0}^2 \right)} + \frac{\lambda_c^4(4/9)}{\left((\hat{s} - m_{D_0^c}^2)^2 + \Gamma_{D_0^c}^2 m_{D_0^c}^2 \right)} \right\}, \quad (3.3)$$

where $\sqrt{\hat{s}}$ is the centre-of-mass energy of the subprocess $qq \rightarrow D_0 \rightarrow qq$, x_1 and x_2 are the fractions of the momenta of the incoming hadrons carried by the quarks with distributions $u(x)$, $d(x)$, etc., and Γ_{D_0, D_0^c} are the decay widths of the D_0 and D_0^c . No mixing between the D_0 and D_0^c is taken into account in writing (3.3) and (3.2): its effects would be trivial.

In fig. 1 we show the total cross sections for the above process as a function of the mass m_{D_0} , assuming $m_{D_0} = m_{D_0^c}$ and $\lambda_Q = \lambda_c \equiv \lambda$ for simplicity. The common Yukawa coupling is parametrized by its ratio to the electromagnetic coupling

$$F \equiv \frac{\lambda^2/4\pi}{\alpha_{em}} = \frac{\lambda_c^2/4\pi}{\alpha_{em}}. \quad (3.4)$$

Note that although neutral current constraints impose severe upper bounds on flavour- or generation-changing couplings of the D supermultiplets, there are no severe constraints on their diagonal couplings [17]. Note also that the three generations of scalar D particles could in principle have similar masses, in which case the estimated cross sections should be increased by a factor of 3. Fig. 1 includes curves for all high-energy hadron-hadron colliders currently envisaged: CERN $\bar{p}p$ at $\sqrt{s} = 630$ GeV, FNAL $\bar{p}p$ at $\sqrt{s} = 1600$ GeV, LHC pp at $\sqrt{s} = 17$ TeV and SSC pp at $\sqrt{s} = 40$ TeV.

These processes will appear experimentally as two-jet events. We have made a detailed comparison of the signal-to-background ratio at $\sqrt{s} = 630$ GeV using data published by UA2 [20]. In table 2 we compare their $d\sigma/dM$ (jet-jet) multiplied by the widths δM of the bins they use, corresponding to their mass resolution, with the

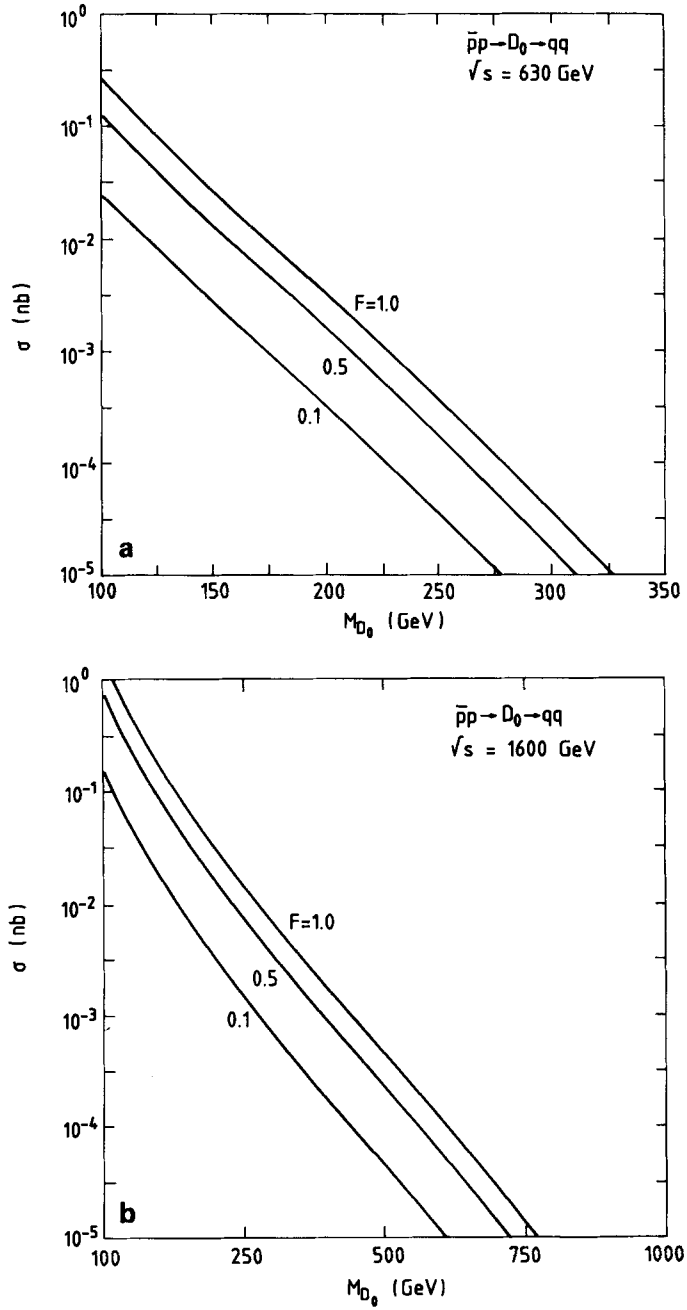


Fig. 1. Total cross sections, as functions of m_{D_0} and $F \equiv (\lambda^2/4\pi)/\alpha_{em}$, for the process $h + h \rightarrow (D \rightarrow qq) + X$ for (a) $\bar{p}p$ at $\sqrt{s} = 630$ GeV, (b) $\bar{p}p$ at $\sqrt{s} = 1600$ GeV, (c) pp at $\sqrt{s} = 17$ TeV and (d) pp at $\sqrt{s} = 40$ TeV. The cross sections include D_0 , D_0^c , \bar{D}_0 and \bar{D}_0^c production.

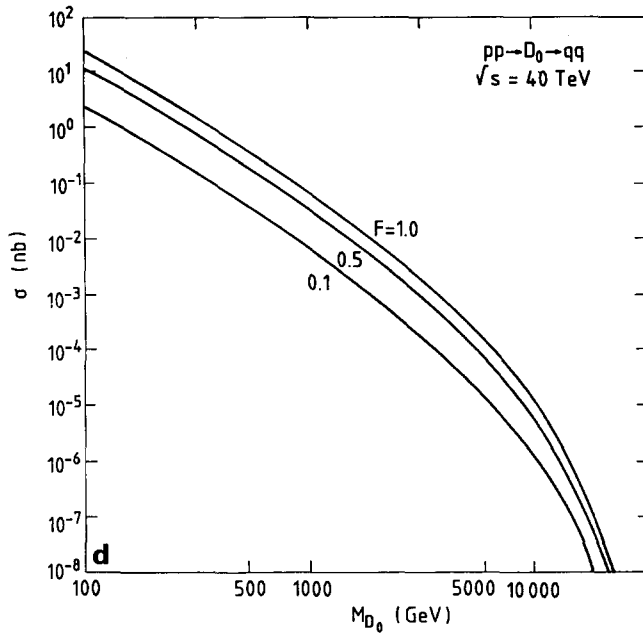
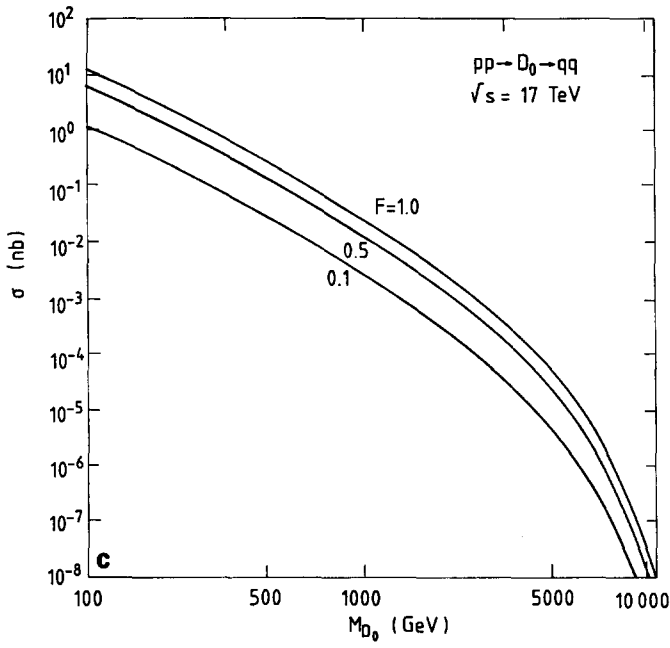


Fig. 1. (continued).

TABLE 2
Signal and background for D_0 search in $\bar{p}p$ collisions

$\sqrt{s} = 630 \text{ GeV}$								
$m \text{ (GeV)}$	125	145	168	188	208	228	268	288
$\delta m \text{ (GeV)}$	11.0	12.3	13.7	14.9	16.2	17.3	19.5	20.6
$d\sigma/dm \text{ (nb/GeV)}$	3×10^{-1}	9×10^{-2}	3×10^{-2}	10^{-2}	7×10^{-3}	3×10^{-3}	4×10^{-4}	6^{-4}
$\delta\sigma \text{ (nb)}$	3.3	1.11	0.41	0.15	0.11	0.05	7.8×10^{-3}	12×10^{-3}
$\sigma \text{ (nb)}$	0.8×10^{-1}	0.3×10^{-1}	0.1×10^{-1}	0.5×10^{-2}	0.2×10^{-2}	0.9×10^{-3}	0.2×10^{-3}	0.6×10^{-4}
$F = 1$								

We have assumed, following UA2 [20], that the mass resolution $\delta m = 0.29 [m \text{ (GeV)}]^{3/4}$ and computed $\delta\sigma = 2\delta m d\sigma/dm$.

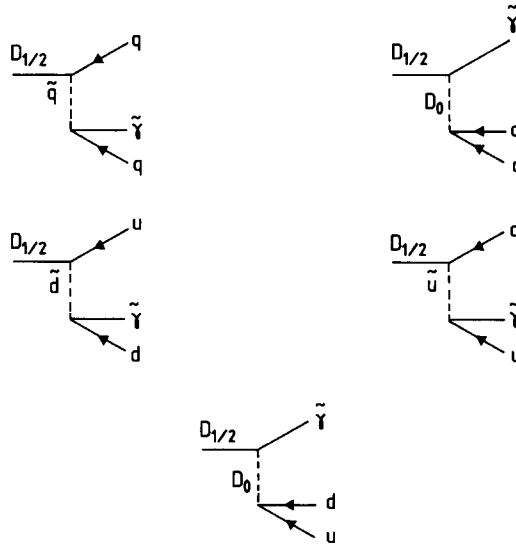


Fig. 2. Tree diagrams contributing to the decay $D_{1/2} \rightarrow \bar{q}\bar{q}\tilde{\chi}$.

D_0 cross section we would expect for $F = 1$. We see that the QCD jet-jet background is between one and two orders of magnitude larger than our expected D_0 cross sections. We expect this state of affairs to be repeated at future colliders, so that it will be difficult to see the D_0 in this way. Nevertheless, it would be useful for theorists if collider experiments could in the future quote directly upper limits on cross sections for new particles decaying into jet pairs.

We turn now to the decays of the fermions $D_{1/2}$. In fig. 2 we show diagrams contributing to their possible decays $D_{1/2} \rightarrow \bar{q} + \bar{q} + \tilde{\chi}$ in leading order. The squared

decay amplitude in the rest frame of the $D_{1/2}$ is given by:

$$\begin{aligned}
 |\mathcal{M}_{D_{1/2}}|^2 &= \lambda_c^2 (\sqrt{2} e)^2 \\
 &\times 4 \left\{ \frac{\left(\frac{1}{3}\right)^2 M_{D_{1/2}} E_1 \left(M_{D_{1/2}} E_1 + \frac{1}{2} M_{D_{1/2}}^2 - \frac{1}{2} m_{\tilde{\gamma}}^2\right)}{\left(P^2 - m_{\tilde{d}_R}^2\right)^2} \right. \\
 &+ \frac{\left(\frac{2}{3}\right)^2 M_{D_{1/2}} E_2 \left(M_{D_{1/2}} E_2 + \frac{1}{2} M_{D_{1/2}}^2 - \frac{1}{2} m_{\tilde{\gamma}}^2\right)}{\left(P^2 - m_{\tilde{u}_R}^2\right)^2} \\
 &+ \frac{\left(\frac{1}{3}\right)^2 M_{D_{1/2}} \left(M_{D_{1/2}} - E_1 - E_2\right) \left(M_{D_{1/2}} (E_1 + E_2) - \frac{1}{2} M_{D_{1/2}}^2 + \frac{1}{2} m_{\tilde{\gamma}}^2\right)}{\left(k^2 - m_{D_0^2}^2\right)^2} \\
 &- \frac{\left(\frac{1}{3}\right) \left(-\frac{2}{3}\right) M_{D_{1/2}}^2 \left(2 E_1 E_2 - M_{D_{1/2}} (E_1 + E_2) + \frac{1}{2} M_{D_{1/2}}^2 - \frac{1}{2} m_{\tilde{\gamma}}^2\right)}{\left(P^2 - m_{\tilde{u}_R}^2\right) \left(P^2 - m_{\tilde{d}_R}^2\right)} \\
 &- \frac{\left(\frac{1}{3}\right) \left(-\frac{1}{3}\right) M_{D_{1/2}} \left(M_{D_{1/2}} - 2 E_2\right) \left(M_{D_{1/2}} (E_1 + E_2) - \frac{1}{2} M_{D_{1/2}}^2 + \frac{1}{2} m_{\tilde{\gamma}}^2\right)}{\left(P^2 - m_{\tilde{d}_R}^2\right) \left(k^2 - m_{D_0^2}^2\right)} \\
 &+ \left. \frac{\left(-\frac{2}{3}\right) \left(-\frac{1}{3}\right) M_{D_{1/2}} \left(M_{D_{1/2}} - 2 E_1\right) \left(M_{D_{1/2}} (E_1 + E_2) - \frac{1}{2} M_{D_{1/2}}^2 + \frac{1}{2} m_{\tilde{\gamma}}^2\right)}{\left(P^2 - m_{\tilde{u}_R}^2\right) \left(k^2 - m_{D_0^2}^2\right)} \right\} \\
 &+ \lambda_Q^2 (\sqrt{2} e)^2 \cdot 4 \left\{ \frac{\left(\frac{1}{3}\right)^2 \cdot 4 \cdot M_{D_{1/2}} \left(M_{D_{1/2}} - E_1 - E_2\right) \left(M_{D_{1/2}} (E_1 + E_2) - \frac{1}{2} M_{D_{1/2}}^2 + \frac{1}{2} m_{\tilde{\gamma}}^2\right)}{\left(k^2 - m_{D_0^2}^2\right)^2} \right. \\
 &+ M_{D_{1/2}} E_1 \left(M_{D_{1/2}} E_1 + \frac{1}{2} M_{D_{1/2}}^2 - \frac{1}{2} m_{\tilde{\gamma}}^2\right) \\
 &\times \left(\frac{\left(\frac{2}{3}\right)^2}{\left(P^2 - m_{\tilde{u}_L}^2\right)^2} + \frac{\left(-\frac{1}{3}\right)^2}{\left(P^2 - m_{\tilde{d}_L}^2\right)^2} + \frac{2 \left(\frac{2}{3}\right) \left(-\frac{1}{3}\right)}{\left(P^2 - m_{\tilde{u}_L}^2\right) \left(P^2 - m_{\tilde{d}_L}^2\right)} \right) \\
 &+ \frac{\left(-\frac{1}{3}\right) M_{D_{1/2}} \left(M_{D_{1/2}} - 2 E_2\right) \left(M_{D_{1/2}} (E_1 + E_2) - \frac{1}{2} M_{D_{1/2}}^2 + \frac{1}{2} m_{\tilde{\gamma}}^2\right)}{\left(k^2 - m_{D_0^2}^2\right)} \\
 &\times \left. \left(\frac{\left(-\frac{2}{3}\right)}{\left(P^2 - m_{\tilde{u}_L}^2\right)} + \frac{\left(\frac{1}{3}\right)}{\left(P^2 - m_{\tilde{d}_L}^2\right)} \right) \right\}, \tag{3.5}
 \end{aligned}$$

where

$$P^2 = M_{D_{1/2}}(M_{D_{1/2}} - 2E_1), \quad k^2 = 2E_1E_2(1 - \cos \phi).$$

ϕ is the angle between the two final-state quarks, E_1 and E_2 their energies, we have assumed that the lightest supersymmetric particle $\tilde{\chi}$ is essentially a photino $\tilde{\gamma}$ [21] and we have neglected mixing between left- and right-handed squarks. The partial decay rate d of the $D_{1/2}$ is given by

$$d\Gamma = \frac{1}{(2\pi)^3} \frac{1}{m_{D_{1/2}}} |\mathcal{M}_{D_{1/2}}|^2 dE_1 dE_2. \quad (3.6)$$

The decay rate given by (3.5) and (3.6) is a complicated function of $m_{D_{1/2}}/m_{\tilde{u},\tilde{d}}$. In the simplifying limit where $m_{D_{1/2}} \ll m_{\tilde{u},\tilde{d}}$ and neglecting quark masses and $m_{\tilde{\chi}}$, we find

$$\Gamma = \frac{1}{(2\pi)^3} M_{D_{1/2}}^5 \cdot \frac{1}{192} \left\{ \lambda_c^2 (\sqrt{2}e)^2 \cdot 4 \left(\frac{(\frac{1}{3})^2}{m_{\tilde{d}_R}^4} + \frac{(-\frac{2}{3})^2}{m_{\tilde{u}_R}^4} + \frac{2(\frac{1}{3})(-\frac{2}{3})}{m_{\tilde{d}_R}^2 m_{\tilde{u}_R}^2} \right) - \lambda_Q^2 (\sqrt{2}e)^2 \cdot 4 \left(\frac{(-\frac{1}{3})^2}{m_{\tilde{d}_L}^4} + \frac{(\frac{2}{3})^2}{m_{\tilde{u}_L}^4} + \frac{2(-\frac{1}{3})(\frac{2}{3})}{m_{\tilde{u}_L}^2 m_{\tilde{d}_L}^2} \right) \right\}. \quad (3.7)$$

Rather than use (3.5) in calculating the missing energy signature for $D_{1/2}$ production discussed in sect. 4, we use simple phase space. We do not expect our results to be sensitive to this simplification.

3.3. COUPLINGS TO QUARKS AND LEPTONS

We now turn to the case (3.1b). Mixing of the D particles with conventional charge $-\frac{1}{3}$ quarks would be possible if $\langle 0|\tilde{p}|0\rangle = 0$. However, as mentioned earlier, this possibility is severely constrained by experiment, and we assume $\langle 0|\tilde{p}|0\rangle = 0$, in which case there is no D/d mixing. Single production of the scalar D is possible in ep collisions via $e^-u \rightarrow D_0/\bar{D}_0^c$ ($e^+\bar{u} \rightarrow D_0^c/\bar{D}_0$), while $D_0/\bar{D}_0^c \rightarrow q + \ell$ ($D_0^c/\bar{D}_0 \rightarrow \bar{q} + \ell$) decays dominate. The squared amplitudes for these decays are

$$|\mathcal{M}_{D_0}|^2 = 6\lambda_e^2 m_{D_0}^2, \quad |\mathcal{M}_{D_0^c}|^2 = 12\lambda_L^2 m_{D_0^c}^2. \quad (3.8)$$

Neglecting mixing, all of the D_0 decays are to jets + charged leptons, while the D_0^c has 50% branching ratios into (jets + charged leptons) and (jets + neutrinos). Again, we have no idea what the ratio of the couplings λ_e and λ_L might be. The cross

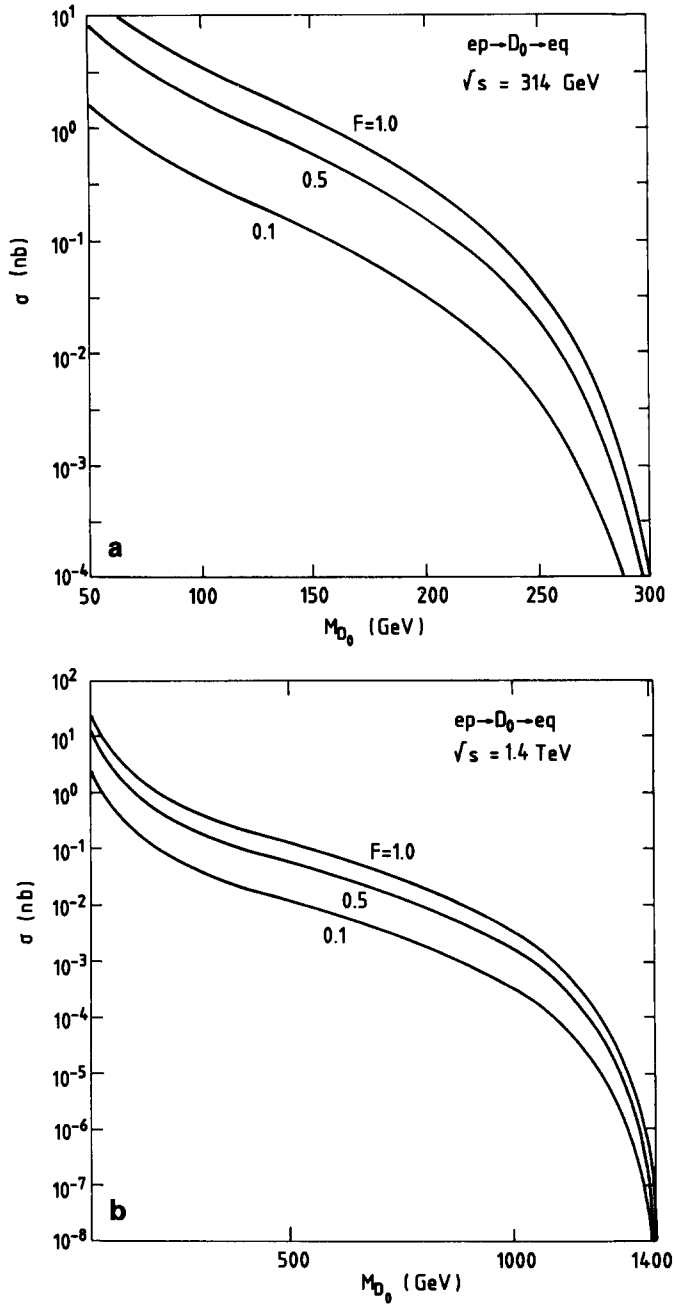


Fig. 3. Total cross section, as a function of m_{D_0} and $F \equiv (\lambda^2/4\pi)/\alpha_{em}$, for the process $ep \rightarrow (D \rightarrow \ell q) + X$; (a) at $\sqrt{s} = 314$ GeV, (b) at $\sqrt{s} = 1.4$ TeV and (c) at $\sqrt{s} = 1.8$ TeV. The cross sections include D_0 and \bar{D}_0^c production.

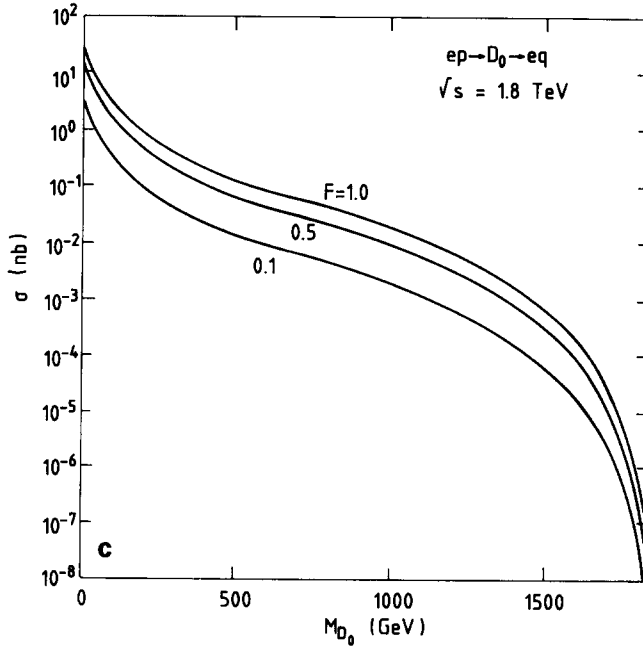


Fig. 3. (continued).

section for the process $ep \rightarrow D_0 \rightarrow \ell q$ is given by

$$d\sigma = \frac{d\hat{s}}{16\pi} \cdot \frac{\hat{s}}{s} \left\{ u \left(\frac{\hat{s}}{s} \right) + \dots \right\} \times \left\{ \frac{2\lambda_L^4}{(\hat{s} - M_{D^c}^2)^2 + \Gamma_{D^c}^2 M_{D^c}^2} + \frac{\lambda_c^4}{(\hat{s} - M_D^2)^2 + \Gamma_D^2 M_D^2} \right\}, \quad (3.9)$$

where $\sqrt{s} = \sqrt{4E_c E_p}$ is the centre-of-mass energy and \hat{s} is the invariant mass squared of the colliding leptoquark combination. In fig. 3 we show the integrated cross section as a function of m_{D_0} , assuming $m_{D_0} = m_{D_0^c}$ and $(\lambda_L^2 + \lambda_c^2) = \lambda^2 = 4\pi \cdot F \cdot \alpha_{em}$ as in eq. (3.4). Results are presented for HERA: $\sqrt{s} = 314$ GeV, and for two LHC/LEP options: $\sqrt{s} = 1.4$ TeV, 1.8 TeV.

These processes will appear experimentally as (lepton + jet) final states. The cases where the lepton is an electron must be compared with the background from the conventional process $e + p \rightarrow e + X$, whereas the cases where the lepton is a neutrino must be compared with the charged current cross sections for $e + p \rightarrow \nu + X$. There would be relatively little background if the D_0 decays into a μ or a τ , but flavour-changing neutral current constraints [17] severely restrict the possible couplings to these leptons.

We show in table 3 a numerical calculations of the cross sections for $e^- + p \rightarrow (D_0 \rightarrow e^- + q) + X$, compared with the continuum background coming from the conventional electromagnetic and neutral current scattering $e^- + p \rightarrow e^- + X$. To estimate the background, we have integrated over bins in x corresponding to the $e + \text{jet}$ mass resolutions expected for the HERA detector ZEUS [22]. We see that the signal-to-background ratios are very favourable if $F = 1$, which may be an upper bound on the $Dq\ell$ coupling. It should also be possible to detect the D_0 even if F is considerably smaller than one. The process $e + q \rightarrow D_0 \rightarrow \nu + q$ can be detected almost equally easily. The background from the standard model charged current reaction is slightly smaller than that shown in table 3. The resolution for the mass bump in the $(\nu + q)$ channel is in the ZEUS detector only a factor ~ 1.6 worse than in the $(e + q)$ channel. The other HERA detector H1 is expected to have an $(e + q)$ mass resolution somewhat better (by a factor ~ 1.1) and a $(\nu + q)$ mass resolution somewhat worse (by a factor $\sim 1/1.7$) than the ZEUS detector. Similar calculations can be made for higher energy ep collider projects, with similar conclusions [23].

The squared amplitude for $D_{1/2} \rightarrow q\ell\tilde{\chi}$ decay in the $D_{1/2}$ rest frame is given by

$$\begin{aligned}
|\mathcal{M}_{D_{1/2}}|^2 &= \lambda_L^2 (\sqrt{2} e)^2 \\
&\times 12 \left\{ \frac{M_{D_{1/2}} E_q \left(\frac{1}{2} M_{D_{1/2}}^2 - \frac{1}{2} m_{\tilde{\gamma}}^2 - M_{D_{1/2}} E_q \right)}{(P^2 - m_{\tilde{e}_L}^2)^2} \right. \\
&\times \left(\frac{1}{2} M_{D_{1/2}}^2 - \frac{1}{2} m_{\tilde{\gamma}}^2 - M_{D_{1/2}} E_\ell \right) M_{D_{1/2}} E_\ell \left[\frac{\left(\frac{2}{3}\right)^2}{(P^2 - m_{\tilde{u}_L}^2)^2} + \frac{\left(\frac{1}{3}\right)^2}{(P^2 - m_{\tilde{d}_L}^2)^2} \right] \\
&+ \frac{2 \cdot \left(\frac{1}{3}\right)^2 M_{D_{1/2}} (M_{D_{1/2}} - E_q - E_\ell) (M_{D_{1/2}} (E_\ell + E_q) - \frac{1}{2} M_{D_{1/2}}^2 + \frac{1}{2} m_{\tilde{\gamma}}^2)}{(k^2 - m_{D_0^*}^2)^2} \\
&+ \frac{\left(\frac{2}{3}\right) \left(2 M_{D_{1/2}}^2 E_\ell E_q - M_{D_{1/2}}^2 (M_{D_{1/2}} (E_\ell + E_q) - \frac{1}{2} M_{D_{1/2}}^2 + \frac{1}{2} m_{\tilde{\gamma}}^2) \right)}{(P^2 - m_{\tilde{u}_L}^2) (P^2 - m_{\tilde{e}_L}^2)} \\
&+ \frac{\left(\frac{1}{3}\right) M_{D_{1/2}} \left(M_{D_{1/2}} (E_\ell + E_q) - \frac{1}{2} M_{D_{1/2}}^2 + \frac{1}{2} m_{\tilde{\gamma}}^2 \right)}{(k^2 - m_{D_0^*}^2)} \\
&\times \left[\frac{M_{D_{1/2}} - 2E_q}{(P^2 - m_{\tilde{e}_L}^2)} + \frac{(M_{D_{1/2}} - 2E_\ell) \left(\frac{1}{3}\right)}{(P^2 - m_{\tilde{d}_L}^2)} + \frac{(M_{D_{1/2}} - 2E_\ell) \left(-\frac{2}{3}\right)}{(P^2 - m_{\tilde{u}_L}^2)} \right] \left. \right\}
\end{aligned}$$

$$\begin{aligned}
& + \lambda_e^2 (\sqrt{2} e)^2 \cdot 12 \left\{ \frac{\left(-\frac{2}{3}\right)^2 M_{D_{1/2}} E_\ell \left(\frac{1}{2} M_{D_{1/2}}^2 - \frac{1}{2} m_{\tilde{\gamma}}^2 - M_{D_{1/2}} E_\ell\right)}{\left(P^2 - m_{\tilde{u}_R}^2\right)^2} \right. \\
& + \frac{M_{D_{1/2}} E_q \left(\frac{1}{2} M_{D_{1/2}}^2 - \frac{1}{2} m_{\tilde{\gamma}}^2 - M_{D_{1/2}} E_q\right)}{\left(P^2 - m_{\tilde{e}_R}^2\right)^2} \\
& - \frac{\left(-\frac{2}{3}\right) \left(2 M_{D_{1/2}}^2 E_\ell E_q - M_{D_{1/2}}^2 \left(M_{D_{1/2}} (E_\ell + E_q) - \frac{1}{2} M_{D_{1/2}}^2 + \frac{1}{2} m_{\tilde{\gamma}}^2\right)\right)}{\left(P^2 - m_{\tilde{e}_R}^2\right) \left(P^2 - m_{\tilde{u}_R}^2\right)} \\
& \left. + \frac{\left(\frac{1}{3}\right)^2 M_{D_{1/2}} \left(M_{D_{1/2}} - E_q - E_\ell\right) \left(M_{D_{1/2}} (E_\ell + E_q) - \frac{1}{2} M_{D_{1/2}}^2 + \frac{1}{2} m_{\tilde{\gamma}}^2\right)}{\left(k^2 - m_{D_0}^2\right)^2} \right\}, \tag{3.10}
\end{aligned}$$

where

$$P^2 = M_{D_{1/2}} \left(M_{D_{1/2}} - 2E_{\ell, q} \right),$$

$$k^2 = 2E_\ell E_q (1 - \cos \omega).$$

ω is the angle between the final-state quark and lepton, E_q and E_ℓ are their

TABLE 3
Signal and background for D_0 search in ep collisions

m_{D_0}	100	200	300
$x = m_{D_0}^2/s$	0.1	0.40	0.91
Background cross sections			
$d\sigma/dx$ e_L^- (nb)	42.3	1.4	2.0×10^{-4}
e_R^- (nb)	24.6	0.7	9.3×10^{-5}
δm (GeV) ZEUS [22]	2.05	2.9	3.5
$\delta\sigma = 4m \delta m\sigma/s$ (nb)	(L) 0.35	0.032	8.4×10^{-6}
(ZEUS) [22]	(R) 0.20	0.016	3.9×10^{-6}
σ_{D_0} ($F=1$) (nb)	3.6	0.32	0.83×10^{-4}

Calculations for HERA: $\sqrt{s} = 314$ GeV, corresponding to $E_e = 30$ GeV, $E_p = 820$ GeV.

energies and $d\Gamma$ is given again by (3.6). In the simplifying limit where $m_{D_{1/2}} \ll m_t, m_{\bar{q}}$, we find

$$\Gamma = \frac{1}{(2\pi)^3} \frac{M_{D_{1/2}}^5}{192} \left\{ \lambda_L^2 (\sqrt{2} e)^2 \cdot 3 \cdot 4 \left(\frac{1}{m_{\bar{c}_L}^4} - \frac{\left(\frac{2}{3}\right)^2}{m_{\bar{u}_L}^4} - \frac{\left(-\frac{1}{3}\right)^2}{m_{\bar{d}_L}^4} - \frac{2\left(\frac{2}{3}\right)^2}{m_{\bar{u}_L}^2 m_{\bar{c}_L}^2} \right) \right. \\ \left. + \lambda_c^2 (\sqrt{2} e)^2 \cdot 3 \cdot 4 \left(\frac{\left(-\frac{2}{3}\right)^2}{m_{\bar{u}_R}^4} + \frac{1}{m_{\bar{e}_R}^4} + \frac{2\left(-\frac{2}{3}\right)^2}{m_{\bar{e}_R}^2 m_{\bar{u}_R}^2} \right) \right\}, \quad (3.11)$$

which is used in calculating signatures for $D_{1/2}$ production in sect. 4.

3.4. COUPLINGS TO CONJUGATE NEUTRINOS

Finally, we turn to the case (3.1c). Mixing of the D particles with conventional charge $-\frac{1}{3}$ quarks would be possible if $\langle 0|\tilde{\nu}^c|0\rangle \neq 0$. However, as mentioned before, we assume here that $\langle 0|\tilde{\nu}^c|0\rangle = 0$ and hence that there is no D/d mixing. In this case there is no single production of the D_0 in either hadron-hadron or electron-hadron collisions. The dominant decays of the D_0 are to $q + \nu^c$, which would have the same experimental missing-energy signature as conventional squark $\tilde{q} \rightarrow q + \tilde{\gamma}$ decay. The dominant decays of the $D_{1/2}$ in this case are to $q + \nu^c + \tilde{\chi}$ via the squared amplitude

$$|\mathcal{M}|^2 = \lambda_\nu^2 (\sqrt{2} e)^2 \cdot 12 \cdot \frac{\left(\frac{1}{3}\right)^2 M_{D_{1/2}} E_\nu \left(\frac{1}{2} M_{D_{1/2}}^2 - \frac{1}{2} m_{\tilde{\gamma}}^2 - M_{D_{1/2}} E_\nu\right)}{\left(P^2 - m_{\bar{u}_R}^2\right)^2}, \\ P^2 = M_{D_{1/2}} (M_{D_{1/2}} - 2E_\nu), \quad (3.12)$$

where we have already taken the simplifying limit $m_{D_{1/2}} \ll m_t, m_{\bar{q}}$ which is used in the phenomenological analysis of its missing energy signature in sect. 4.

4. Pair production in hadron-hadron collisions

In this section we discuss the possible cross sections and signatures for $(\bar{p})p \rightarrow D_0 \bar{D}_0 + X$ and $D_{1/2} \bar{D}_{1/2} + X$.

4.1. CROSS SECTIONS

The forms of the parton-parton cross sections for $gg, \bar{q}q \rightarrow D_0 \bar{D}_0$ are identical to those of $gg, \bar{q}q \rightarrow \tilde{q}\tilde{q}$ [24] if one compares the limits $m_{D_{1/2}} \gg m_{D_0}$ and $m_{\tilde{g}} \gg m_{\bar{q}}$,

due to the fact that D_0 couplings to gluons are identical to those of \tilde{q} . The only difference between the total cross sections in these limits is therefore an overall combinatorial factor counting the total number of D_0 or \tilde{q} species. In previous work, two of us [25] have generally assumed 5 approximately degenerate flavours of squark (\tilde{u} , \tilde{d} , \tilde{s} , \tilde{c} and \tilde{b}) and added together both left- and right-handed quarks. Here we add together the D_0 scalars expected from three generations, and also include both left- and right-handed states. In this case, $\sigma(D_0\bar{D}_0)/\sigma(\tilde{q}\tilde{\bar{q}}) = \frac{3}{5}$.

The forms of the parton-parton cross sections for $gg, \bar{q}q \rightarrow D_{1/2}\bar{D}_{1/2}$ are identical to those of $gg, \bar{q}q \rightarrow t\bar{t}$ [26] if one considers the limit $m_{D_0} \gg m_{D_{1/2}}$. Equivalently, these parton-parton cross sections can be obtained from those for $gg, \bar{q}q \rightarrow \tilde{g}\tilde{g}$ [24] in the limit $m_{\tilde{q}} \gg m_{\tilde{g}}$ simply by adjusting the colour factors.

4.2. POSSIBLE SIGNATURES

As discussed in the previous section, possible decays of the D_0 are to $q\nu$, $q\ell^-$ and $\bar{q}\bar{q}$, while the $D_{1/2}$ may decay to $q\nu\tilde{\chi}$, $q\ell^-\tilde{\chi}$ or $\bar{q}\bar{q}\tilde{\chi}$, where the $\tilde{\chi}$ is a weakly-interacting neutral Majorana fermion similar to the photino, which can carry off missing energy. We therefore have the following possible event signatures:

missing energy:	$D_0\bar{D}_0 \rightarrow (q\nu)(\bar{q}\bar{\nu})$
	$D_{1/2}\bar{D}_{1/2} \rightarrow (\bar{q}\bar{q}\tilde{\chi})(qq\tilde{\chi})$ or $(q\nu\tilde{\chi})(\bar{q}\bar{\nu}\tilde{\chi})$,
charged lepton pairs:	$D_0\bar{D}_0 \rightarrow (q\ell^-)(\bar{q}\ell^+)$,
leptons and missing energy:	$D_{1/2}\bar{D}_{1/2} \rightarrow (q\ell^-\tilde{\chi})(\bar{q}\ell^+\tilde{\chi})$,
dijet mass bumps:	$D_0\bar{D}_0 \rightarrow (\bar{q}\bar{q})(qq)$.

While the semileptonic decay of one D is not compatible with the simultaneous hadronic decay of another D , it is in principle possible to mix the semileptonic decay into a charged lepton of one D with the neutrino decay of another D . However, we will not discuss such dijet + lepton + missing energy final states. Nor will we discuss final states with two charged leptons of different flavours: $e^\pm\mu^\mp$, $e^\pm\pi^\mp$, $\mu^\pm\tau^\mp$. We concentrate on the missing-energy signatures and on the charged lepton pair signatures. We present here results for $\bar{p}p$ collisions at $\sqrt{s} = 630$ GeV corresponding to the CERN Collider, and 1600 GeV, corresponding to the Fermilab Collider: results for pp collisions at $\sqrt{s} = 17$ TeV corresponding to the LHC will be presented elsewhere [23].

4.3. MISSING ENERGY SIGNATURES

To discuss these quantitatively we have used the same approach as in previous work [25] on UA1 data [14]. We divide missing energy events into the following categories; monojets: only one cluster of hadronic energy within $\Delta R = \sqrt{(\Delta\phi)^2 + (\Delta\eta)^2} = 1$ above a threshold of $E_T = 12$ GeV, and missing transverse

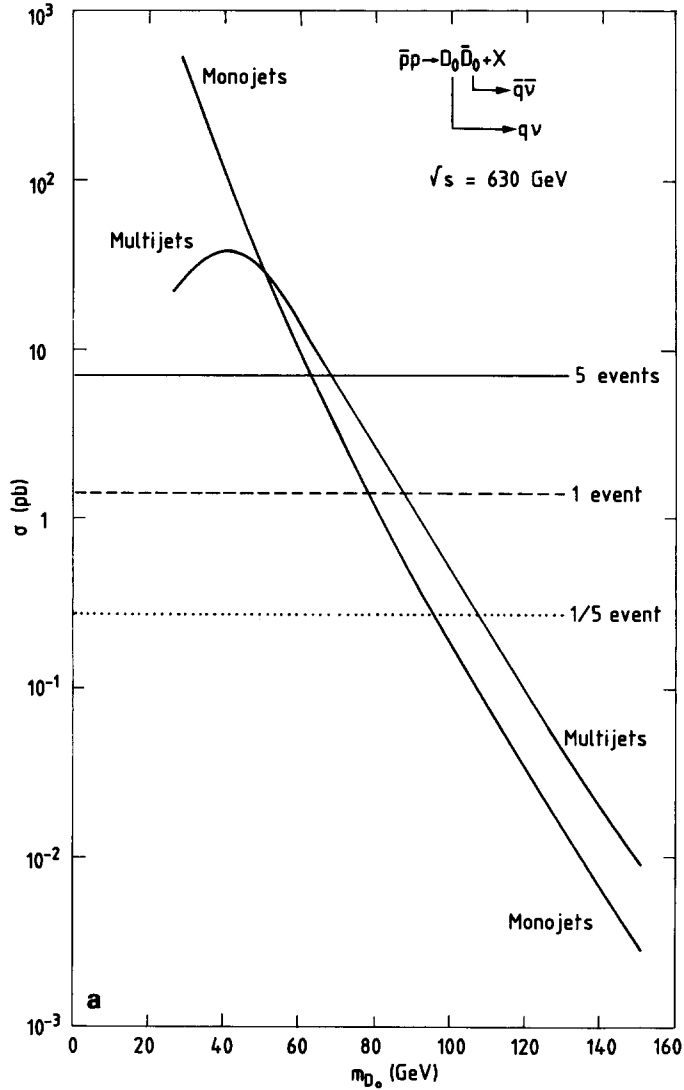


Fig. 4. Total cross section, as a function of m_{D_0} , for the process $\bar{p}p \rightarrow (D_0 \rightarrow q\nu) + (\bar{D}_0 \rightarrow \bar{q}\bar{\nu}) + X$, $(D_0^c \rightarrow \bar{q}\bar{\nu}) + (\bar{D}_0^c \rightarrow q\nu) + X$; (a) at $\sqrt{s} = 630$ GeV and (b) at $\sqrt{s} = 1600$ GeV.

momentum p_T in excess of 4σ , where the measurement error $\sigma = 0.7\sqrt{E_T}$ (GeV); dijets: two such jet clusters and $p_T > 4\sigma$; trijets: three such jets, etc. Our calculations include a somewhat more subtle characterization of the UA1 detector which is described in previous publications [25]. However, a full description of the assignments of missing energy events to different categories is impossible in the absence of a full detector simulation. Therefore in this paper we restrict ourselves to quoting

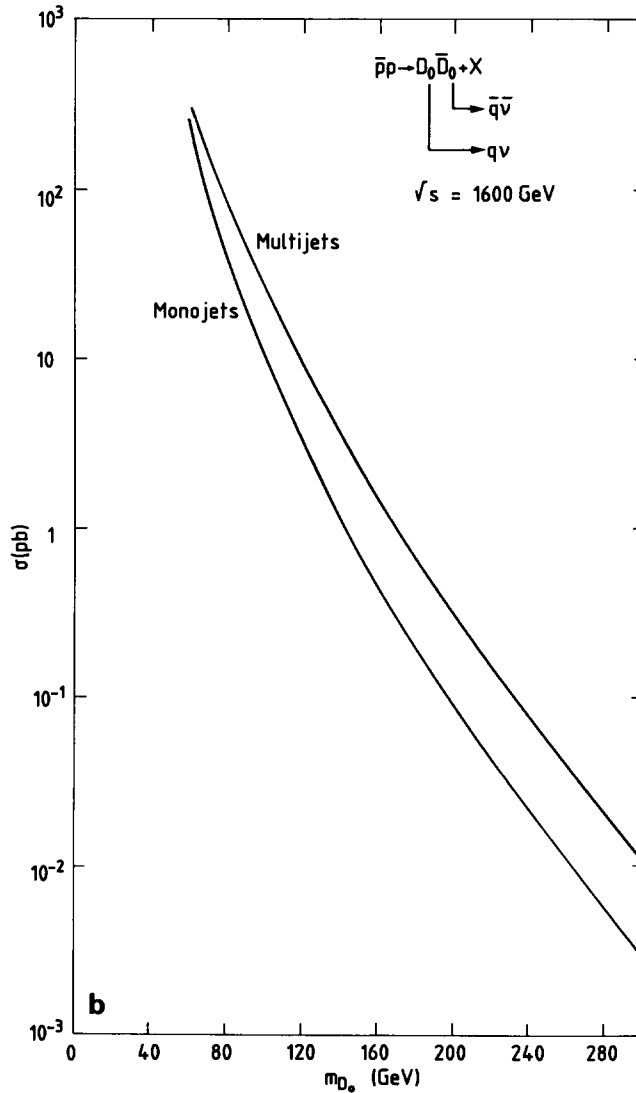


Fig. 4. (continued).

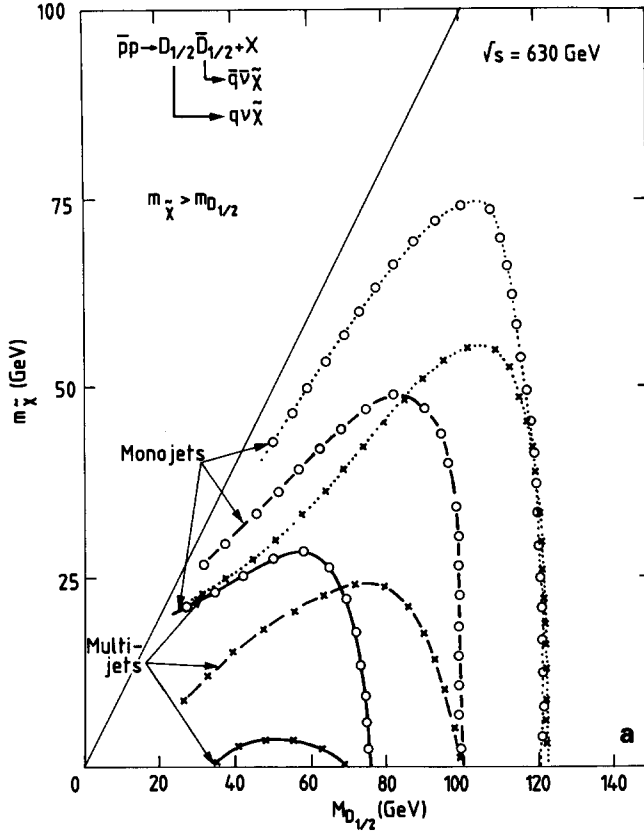


Fig. 5. (a) Contours, in the $(M_{D_{1/2}}, m_{\tilde{\chi}})$ plane, of $\sigma = 7$ pb (solid line), $\sigma = 1.4$ pb (dashed line) and $\sigma = 0.28$ pb (dotted line), for monojet (circles) and multijet (crosses) events from the process $\bar{p}p \rightarrow (D_{1/2} \rightarrow q\nu\tilde{\chi}) + (\bar{D}_{1/2} \rightarrow \bar{q}\bar{\nu}\tilde{\chi}) + X$ at $\sqrt{s} = 630$ GeV; (b) total cross section as a function of $m_{D_{1/2}}$ (assumed equal to $2m_{\tilde{\chi}}$), for the same process at $\sqrt{s} = 1600$ GeV.

cross sections for monojet events and multijet events, and emphasize that a realistic detector may shuffle events between these two categories.

$D_0 \rightarrow q\nu$. The signature for this decay is identical to that previously discussed [13] for $\tilde{q} \rightarrow q\tilde{\gamma}$ with $m_{\tilde{\gamma}} = 0$. Fig. 4a contains our predictions for monojets and multijets* from $D_0\bar{D}_0$ production at $\sqrt{s} = 630$ GeV including the factor of $\frac{3}{5}$ in the cross section which was mentioned previously. The horizontal lines correspond to $\sigma = 7$ pb (solid line: five events in the present event sample of about 700 nb^{-1}), $\sigma = 1.4$ pb (dashed line: corresponding to one event in the present sample) and $\sigma = 0.28 \text{ pb}^{-1}$ (dotted line: corresponding to perhaps five events at the $\bar{p}p$ collider with ACO [27]). We believe that a lower limit on the D_0 mass could only be

* Note that in this and the next case the vast majority of multijet events only contain two jets.

established by the UA1 collaboration itself. If it would establish an upper limit of five multijet events in the present data, that would correspond to $m_{D_0} \geq 60$ to 70 GeV. Fig. 4b contains analogous cross section curves for $\sqrt{s} = 1600$ GeV. At this energy, a sensitivity comparable to present CERN Collider data would increase the possible limit to $m_{D_0} \geq 120$ to 130 GeV.

$D_{1/2} \rightarrow q\nu\tilde{\chi}$. Here there is some additional ambiguity provided by the unknown mass of the $\tilde{\chi}$. Model analyses of sparticle spectra suggest that $m_{\tilde{\chi}} \geq 15$ GeV, and it

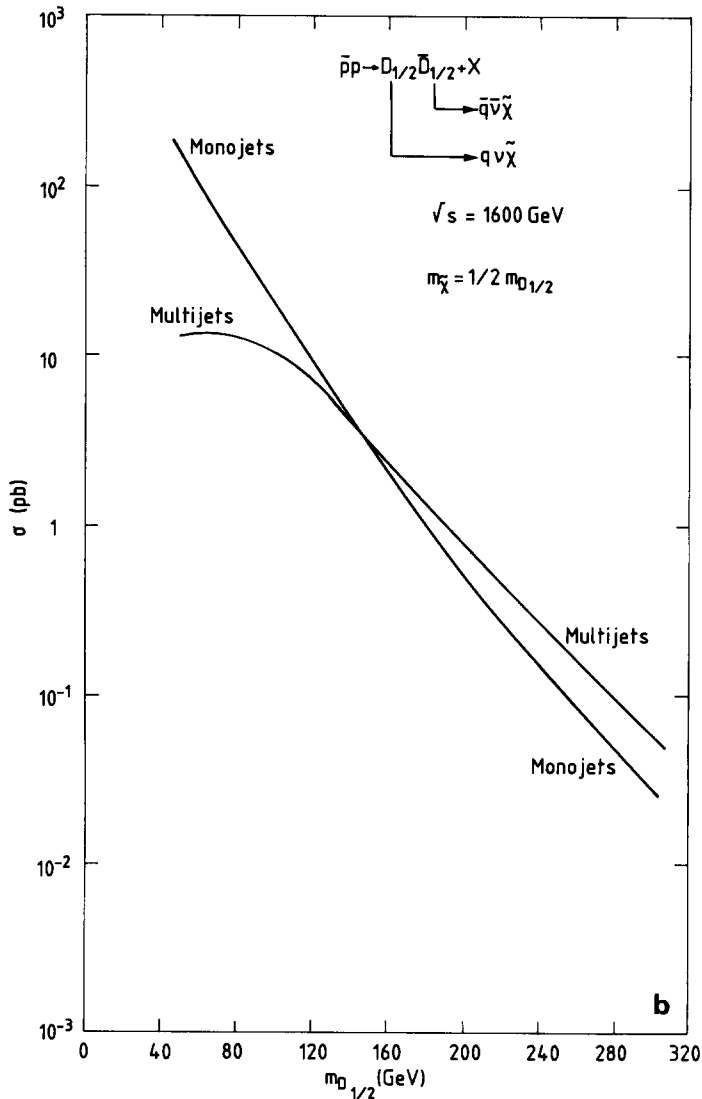


Fig. 5. (continued).

could easily be a large fraction of $m_{D_{1/2}}$. Therefore we plot in fig. 5a our results for $\sqrt{s} = 630$ GeV as contours in the $(m_{D_{1/2}}, m_{\tilde{\chi}}$) plane corresponding to $\sigma = 7$ pb (solid lines), 1.4 pb (dashed lines) and 0.28 pb (dotted lines) as in fig. 4a. Naïve interpolation between these contours will give cross sections accurate to better than a factor of 2, which is in any case the expected accuracy of our calculations. We see that monojet events are always more copious than multijet events. We also see from fig. 5 that the UA1 sensitivity to $m_{D_{1/2}}$ is essentially unchanged for $0 \leq m_{\tilde{\chi}} \leq 15$ GeV, but decreases significantly for larger $m_{\tilde{\chi}}$, and disappears in the limit $m_{\tilde{\chi}} \rightarrow m_{D_{1/2}}$. Less than five monojet events in the present UA1 data [14] would correspond to $m_{D_{1/2}} \geq 70$ GeV if the $\tilde{\chi}$ is light, or $m_{D_{1/2}} \geq 50$ to 60 GeV if $m_{\tilde{\chi}} = \frac{1}{2}m_{D_{1/2}}$. We plot in fig. 5b cross section curves for $\sqrt{s} = 1600$ GeV assuming for definiteness that $m_{\tilde{\chi}} = \frac{1}{2}m_{D_{1/2}}$. We see that a sensitivity comparable to present UA1 data would increase the possible limit to $m_{D_{1/2}} \geq 120$ GeV.

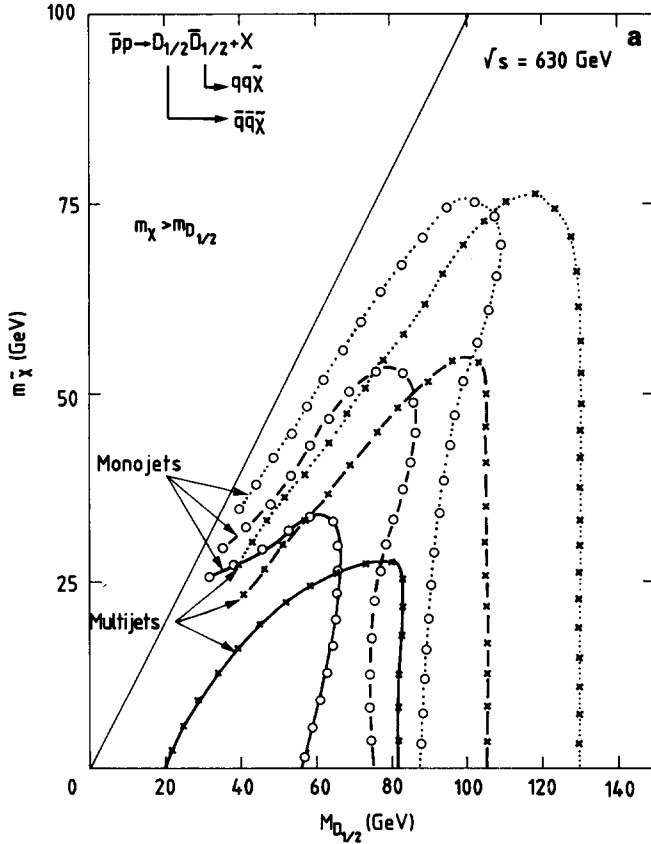


Fig. 6. (a) As in fig. 5a for the process $\bar{p}p \rightarrow (D_{1/2} \rightarrow \bar{q}q\tilde{\chi}) + (\bar{D}_{1/2} \rightarrow qq\tilde{\chi}) + X$ at $\sqrt{s} = 630$ GeV; (b) as in fig. 5b for the same process.

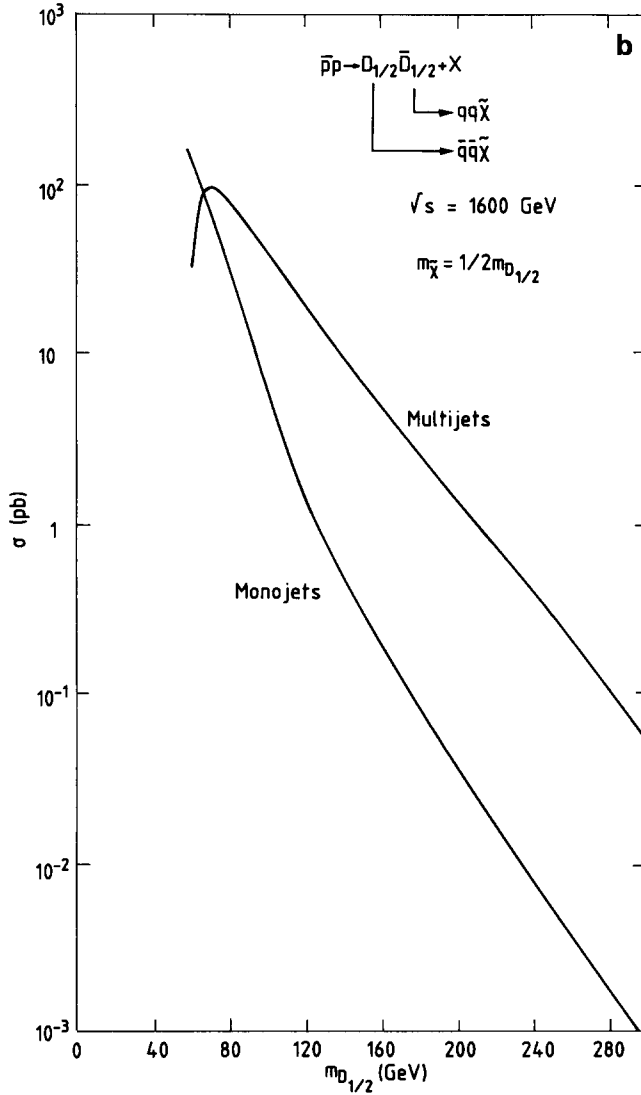


Fig. 6. (continued).

$D_{1/2} \rightarrow \bar{q}q\tilde{\chi}$. This is an alternative decay mode of the $D_{1/2}$ which has a signature similar to the conventional $\tilde{g} \rightarrow q\bar{q}\tilde{\gamma}$ decay, although the cross section is somewhat different and the mass of the $\tilde{\chi}$ may not be negligible. In fig. 6a we have plotted for $\sqrt{s} = 630 \text{ GeV}$ contours in the $(m_{D_{1/2}}, m_{\tilde{\chi}})$ plane corresponding to $\sigma = 7 \text{ pb}$ (solid lines), 1.4 pb (dashed lines) and 0.28 pb (dotted lines) as in fig. 5a. We see that the multijet cross sections are larger for low $m_{\tilde{\chi}}$ and large $m_{D_{1/2}}$, while the multijet cross

sections are larger for bigger $m_{\tilde{\chi}}$. Thus monojet and multijet searches to some extent complement each other. An upper limit of five multijet events in the present data would correspond to $m_{D_{1/2}} \geq 80$ GeV for $m_{\tilde{\chi}} \leq 20$ GeV, while an upper limit of five monojet events in the present data would correspond to $m_{D_{1/2}} \geq 60$ GeV if $m_{\tilde{\chi}} = \frac{1}{2}m_{D_{1/2}}$. Cross section curves for $\sqrt{s} = 1600$ GeV and $m_{\tilde{\chi}} = \frac{1}{2}m_{D_{1/2}}$ are plotted in fig. 6b. Here we see that a multijet sensitivity comparable to that presently achieved by UA1 would reach $m_{D_{1/2}} \sim 150$ GeV.

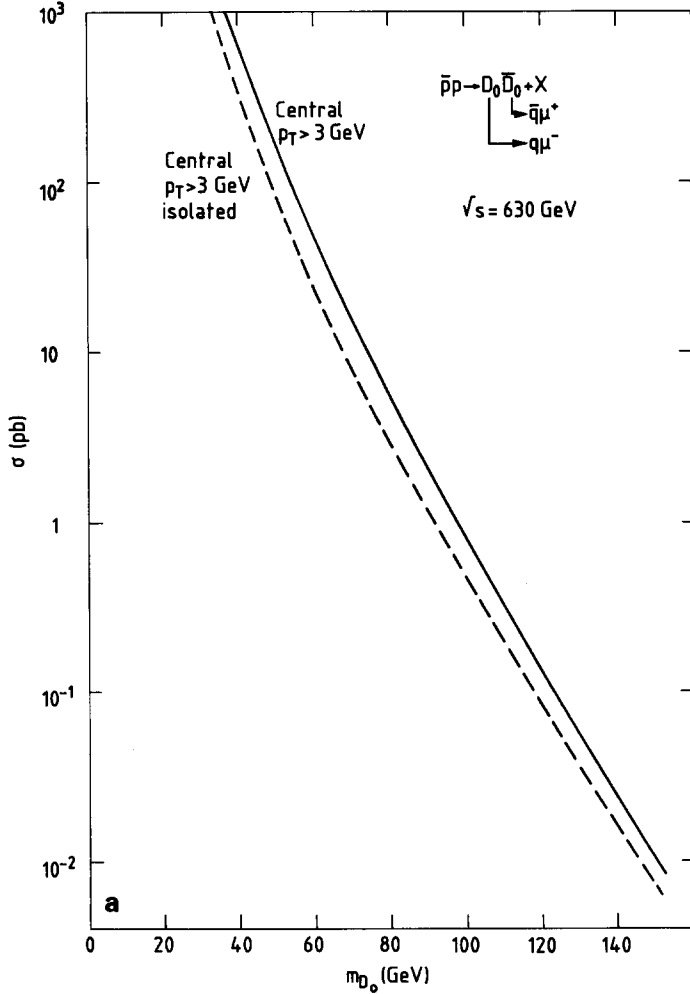


Fig. 7. Total cross section as a function of m_{D_0} for the process $\bar{p}p \rightarrow (D_0 \rightarrow q\mu^-) + (\bar{D}_0 \rightarrow \bar{q}\mu^+) + X$, $(D_0^c \rightarrow \bar{q}\mu^+) + (\bar{D}_0^c \rightarrow q\mu^-) + X$; (a) at $\sqrt{s} = 630$ GeV and (b) at $\sqrt{s} = 1600$ GeV. The centrality and isolation cuts are specified in the text.

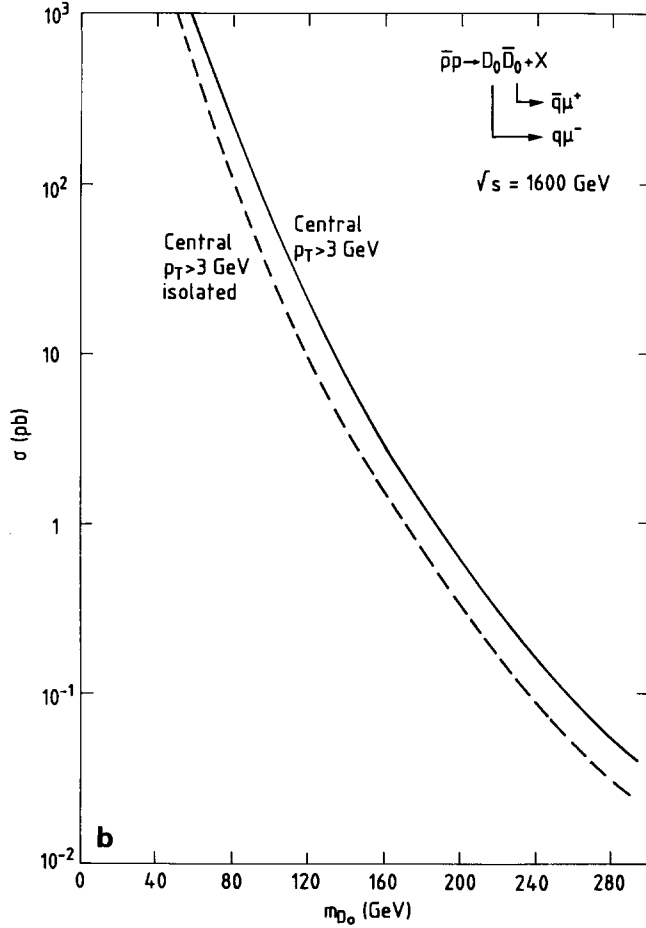


Fig. 7. (continued).

4.4. CHARGED LEPTON SIGNATURES

Taking our cue again from UA1 [15], we have taken the following cuts on charged leptons: $|\eta_1| < 1.3$, $|\eta_2| < 2.0$, $m_{\ell^+\ell^-} > 6 \text{ GeV}$ and $p_T^{\ell_1}, p_T^{\ell_2} > 3 \text{ GeV}$ for muons, $p_T^{\ell_1}, p_T^{\ell_2} > 8 \text{ GeV}$ for electrons. We have also tried the effects of isolation cuts: $(\sum_{1,2} E_T^2) < 9 \text{ GeV}^2$ in the combination of cores with $\Delta R \equiv \sqrt{(\Delta\phi)^2 + (\Delta\eta)^2} < 0.7$ around the two charged leptons. Thus we quote cross sections for $(\mu^+\mu^-)$ both total and isolated, and similarly for (e^+e^-) pairs.

$D_0 \rightarrow q\ell$. Our results at $\sqrt{s} = 630 \text{ GeV}$ are shown in fig. 7a. The total (solid line) and isolated (dashed line) curves are for $\mu^+\mu^-$ pairs. The e^+e^- cross sections are indistinguishable for $m_{D_0} \geq 60 \text{ GeV}$ and differ by less than 10% even for $m_{D_0} = 40 \text{ GeV}$. Note that here we have assumed a 100% branching ratio into $q\mu^-$ for each of

the three generations of D_0 particles. This is probably unreasonable, a better guess being that at most one of the three generations of D_0 particles would have a large branching ratio into $q\mu^-$. In this case, the cross sections in fig. 7 should be reduced by a factor of 3. Taking this point of view, an upper limit of five events in the present UA1 event sample [15] would correspond to $m_{D_0} \geq 60$ to 70 GeV. Corresponding cross sections for $\sqrt{s} = 1600$ GeV are shown in fig. 7b. A sensitivity comparable to UA1's present achievement would then yield $m_{D_0} \geq 140$ GeV.

$D_{1/2} \rightarrow q\ell\tilde{\chi}$. Here we meet again the ambiguity in the mass of the $\tilde{\chi}$. If these events are treated simply as possible missing energy events, and no attempt is made to identify the charged leptons, then the results of fig. 6 are directly applicable.

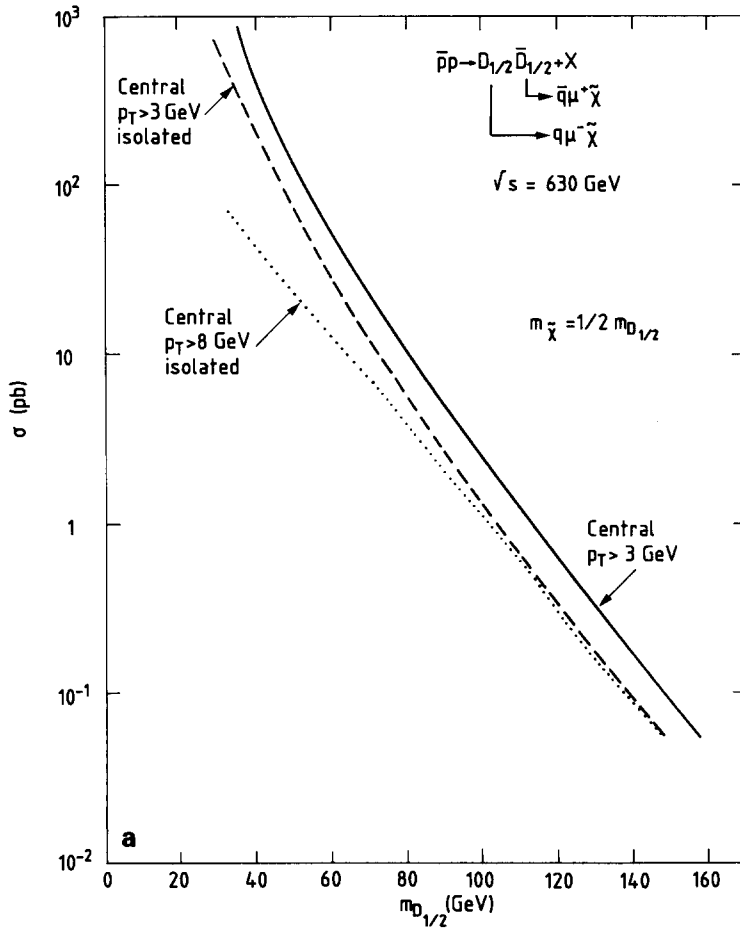


Fig. 8. Same as in fig. 7 for the process $\bar{p}p \rightarrow (D_{1/2} \rightarrow q\mu^- \tilde{\chi}) + (\bar{D}_{1/2} \rightarrow \bar{q}\mu^+ \tilde{\chi}) + X$; (a) at $\sqrt{s} = 630$ GeV and (b) at $\sqrt{s} = 1600$ GeV.

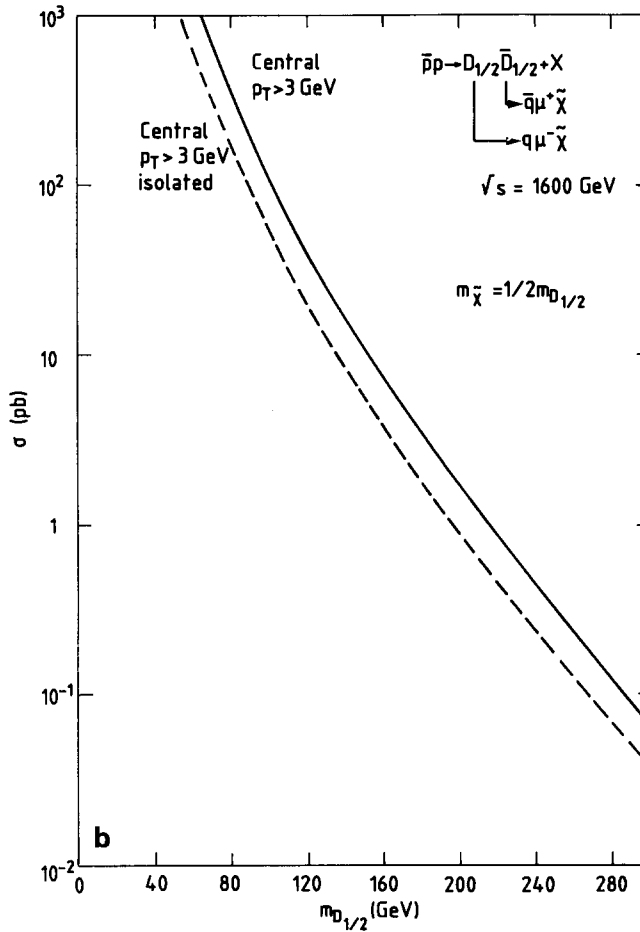


Fig. 8. (continued).

Alternatively, they could be analyzed as $(\ell^+ \ell^-)$ events with no attempt made to measure missing energy. In this case, their signature would be similar to that of $D_0 \rightarrow q\ell$. Figs. 8a and 8b show cross section curves at $\sqrt{s} = 630$ GeV and 1600 GeV respectively, assuming $m_{\tilde{\chi}} = \frac{1}{2}m_{D_{1/2}}$. We see from fig. 8a that an upper bound of five events in the present UA1 sample would yield $m_{D_{1/2}} \geq 80$ GeV, while comparable sensitivity at the Fermilab Collider would yield $m_{D_{1/2}} \geq 160$ GeV.

5. Conclusions

Superstring-inspired models offer the possibility that additional light colour-triplet, charge $|\frac{1}{3}|$ particles exist in addition to conventional quarks. These D

particles exist with spin zero (two per generation) and spin $\frac{1}{2}$ (one Dirac fermion per generation). We have motivated and studied the case where the D particles do not mix with the conventional charge $|\frac{1}{3}|$ quarks d, s, b. In contrast to conventional GUTs, symmetry breaking by the Hosotani mechanism at the compactification scale offers the possibilities that D particles have *either* purely leptoquark couplings $D\bar{q}\bar{\ell}$, or purely diquark couplings Dqq . In this paper we have studied the phenomenologies of these two cases.

In the case of leptoquark couplings, we find that the spin-zero D_0 particles could be produced and observed at HERA if they have masses ≤ 250 GeV. Future ep colliders could extend the search for leptoquark D_0 particles above 1 TeV.

In the case of diquark couplings, we find that the single production of spin-zero D_0 particles at hadron-hadron colliders through qq annihilation is likely to be overwhelmed by a large QCD jet-jet background.

Pair production of $D_0\bar{D}_0$ and $D_{1/2}\bar{D}_{1/2}$ at hadron-hadron colliders offers a promising way to search for these particles. The pair-production mechanism leads to characteristic experimental signatures in almost all possible decay channels if $m_{\tilde{\chi}}/m_D \leq \frac{1}{2}$. The D_0 and $D_{1/2}$ particles could be sought in samples of missing-energy events, $\ell^+\ell^-$ events and top candidates. The present CERN $p\bar{p}$ Collider probably has sensitivity to these particles if they have masses ≤ 60 to 70 GeV. The improved luminosity available with ACOL would increase this range to about 100 GeV. The FNAL Tevatron Collider should be able to see D_0 or $D_{1/2}$ particles weighing up to ~ 150 GeV. In the longer term, at the LHC or the SSC one could search for the pair production of D_0 and $D_{1/2}$ particles with masses ≤ 2 TeV, as has previously been discussed for squark and gluino searches.

We hope that this paper stimulates our experimental colleagues to search systematically in their present and future data for exotic charge $|\frac{1}{3}|$ particles with diquark or leptoquark couplings.

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