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NON-STANDARD HIGGS BOSONS AS
A SOURCE OF MONOJETS

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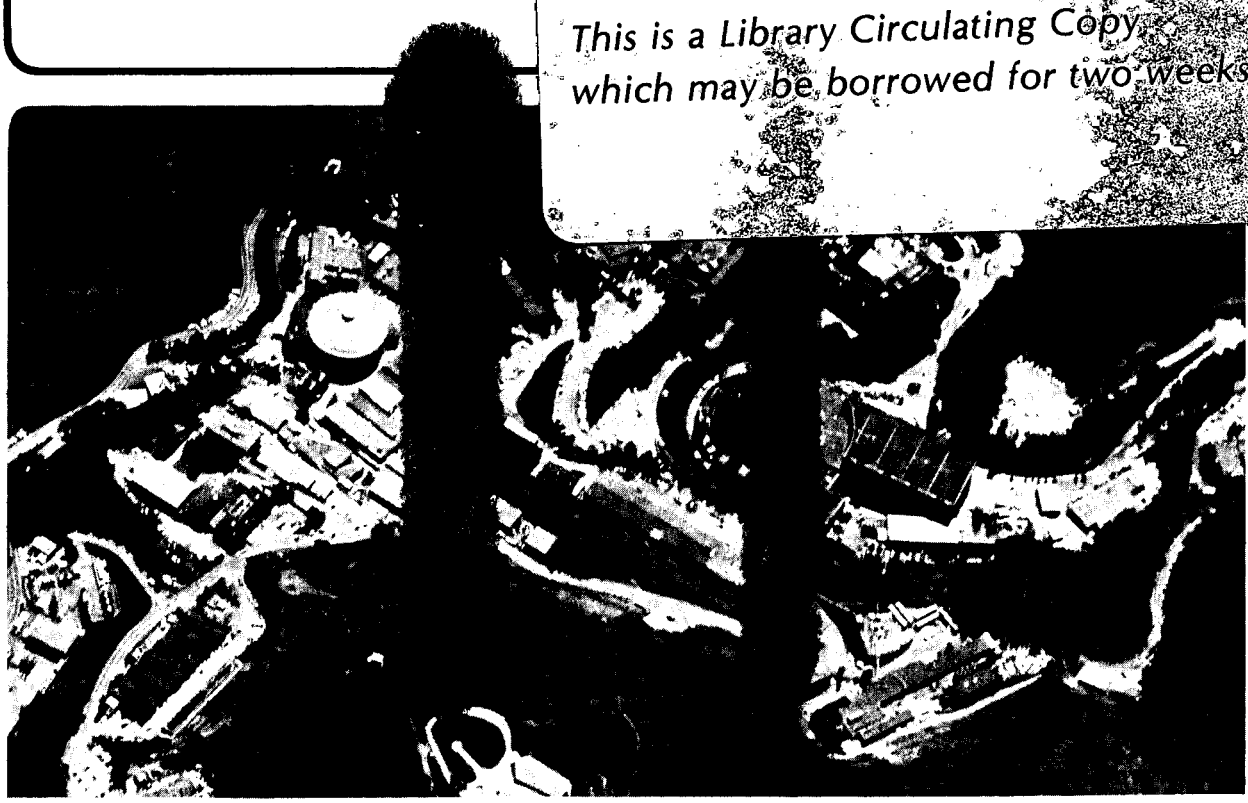

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NON-STANDARD HIGGS BOSONS AS A SOURCE OF MONOJETS

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Abstract

Non-standard Higgs models are examined as a potential source for events with large missing p_T of the type recently seen by the UA1 group at CERN.

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I. INTRODUCTION

The recent discovery of anomalous "monojet" events at the UA1 detector at CERN[1] has generated an increasing amount of excitement among theorists. These events consist of one or more jets and large ($\geq$ 30 GeV) missing transverse energy and are difficult to explain in the context of the standard QCD plus Weinberg-Salam model.[2] In the spring of 1983, UA1 found six such events in a data sample with an integrated luminosity of 113 nb^{-1}.[1]

A number of features make these events interesting and pose problems for model builders. The first is the relatively large rates at which these events occur. These events are also unusual in that the jets look different from typical QCD jets - the monojets have low effective masses and low charged track multiplicities.[1] Also unusual is the comparatively flat missing transverse energy spectra of these events. In this paper, we analyze the restrictions which these facts place upon any attempts made to explain the monojets. There are many explanations in the literature[2] (supersymmetry, strongly interacting W's, etc...) which explain the qualitative features of the data, but some of these run into difficulties when attempting to explain the detailed structure of the monojets. The most natural explanation is in terms of supersymmetry which has classic missing p_T signatures.[3] In this paper, we will discuss some alternatives.

We have not found any elegant model which explains all of the features of the data. In Section II, we describe a model in which there exists more than one Higgs doublet. The monojets then result from the resonant production of a charged Higgs boson. It is clear that the

couplings of this boson to fermions must be extremely non-standard in order to explain the monojet production rate. Our model, however, explains the flat p_T spectra, the low multiplicity, and the low effective mass of the monojets. In the remainder of this introduction, we describe the chain of logic which led us to our model.

What are the possible sources of this missing transverse momentum, p_T^m? Suppose that new particles are produced in partonic collisions. Missing transverse momenta can arise when one of these particles, or one of their decay products, does not interact inside the detector.

The simplest example involves the production of the W boson which then decays to $e\nu$. The p_T^m spectrum caused by the neutrino has a Jacobian peak close to $M_W/2$. If we wish to generate the monojets by a mechanism similar to this, we must produce a particle (Y) as a resonance in parton-parton scattering and then have it decay to a jet plus neutrino(s). The object must have mass at least as large as twice the maximum value of p_T^m observed. So we will require an object of mass at least 150 GeV. In order to generate enough events Y must couple to quarks with at least electroweak strength.

An alternative involves the production of a pair of new particles which decay into objects which carry off unobserved momenta. The most fashionable example of this type is supersymmetry where photinos carry off the missing momentum. In this case, the p_T^m spectrum will begin at some value p_T^0, corresponding to the threshold of the new particle production and will then fall rapidly with increasing p_T^m. (For a specific calculation see Barnett et al.[3].) In this case, the observed rate of monojets requires that the new particles be produced

with rates typical of strong interaction processes. This type of mechanism has been studied in great detail and we do not consider it further.

In order to be useful, the first mechanism must be modified so as to flatten out the Jacobian peak in the p_T^m spectrum. The simplest way for this to occur is to make the resonance (Y) decay into X + neutrino and then force X to decay into jet(s) + neutrino(s). The transverse momenta carried off by the neutrinos tend to cancel and the peak is smeared out.

As was mentioned above, the multiplicity and the invariant mass of the monojets are rather low; less than that of typical jets of the same transverse momenta at the CERN collider. The explanation of the monojet events in terms of supersymmetry produces events which may have too large a multiplicity.[4] This difficulty can be circumvented in the first mechanism if the particle X has a rather low mass since the mass scale M which determines the multiplicity in this case is the mass of X (at least if Y is not strongly interacting).

The least radical scenario which fits these requirements is for Y to be a charged Higgs boson and for X to be the tau lepton or a new heavy lepton. (The standard source of taus, $W \rightarrow \tau\bar{\nu}$, has $p_T^m < M_W/2$ and so cannot explain all the data.) Section II presents a model of this type and Section III contains the numerical results of our study - we expect these results to be applicable to a wide class of models. In Section IV, we present our conclusions.

II. MULTIPLE HIGGS MODELS

We begin by considering a model with two SU(2) Higgs doublets, $\tilde{H}_\alpha$,[5]

$$\tilde{H}_1 = \begin{bmatrix} \tilde{H}_1^{\,0} \\ \tilde{H}_1^{\,-} \end{bmatrix} \quad , \tilde{H}_2 = \begin{bmatrix} \tilde{H}_2^{\,0} \\ \tilde{H}_2^{\,-} \end{bmatrix} \tag{2.1}$$

The vacuum expectation values of these doublets are,

$$<\tilde{H}_1> = \begin{bmatrix} v_1 \\ 0 \end{bmatrix} , <\tilde{H}_2> = \begin{bmatrix} v_2 \\ 0 \end{bmatrix} \tag{2.2}$$

To find the physical particles it is convenient to rotate the fields,

$$\begin{aligned} H_1 &= \frac{1}{\sqrt{v_1^{\,2} + v_2^{\,2}}} \{v_1 \tilde{H}_1 + v_2 \tilde{H}_2\} \\ H_2 &= \frac{1}{\sqrt{v_1^{\,2} + v_2^{\,2}}} \{ - v_2 \tilde{H}_1 + v_1 \tilde{H}_2\} \end{aligned} \tag{2.3}$$

Now the vacuum expectation values are,

$$\begin{aligned} <H_1> &= \sqrt{v_1^{\,2} + v_2^{\,2}} = \frac{gM_w}{2} \\ <H_2> &= 0 \end{aligned} \tag{2.4}$$

The charged components of H_1 and Im $H_1^{\,0}$ become the longitudinal components of the $W^\pm$ and Z bosons. The physical Higgs bosons are $H_2^{\,\pm}$, two linear combinations of Re $H_1^{\,0}$ and Re $H_2^{\,0}$ and the pseudoscalar boson Im $H_2^{\,0}$.

We consider two possible couplings of the fermions to the Higgs bosons which illustrate the problems faced when constructing models of this type. In the first model, the charge $\frac{2}{3}$ quarks get their mass from couplings to $\tilde{H}_1$ and the charge $-\frac{1}{3}$ quarks get their mass from couplings to $\tilde{H}_2$. The Lagrangian coupling the charged Higgs to the fermions is,

$$\pounds = \frac{g}{\sqrt{2M_w}} \left[\frac{v_2}{v_1} \bar{U}_R M_u K D_L + \frac{v_1}{v_2} \bar{U}_L M_d K D_R \right] H_2^{\,+} + \text{h.c.} \tag{2.5}$$

where

$$U = \begin{bmatrix} u \\ c \\ t \end{bmatrix} \quad D = \begin{bmatrix} d \\ s \\ b \end{bmatrix} \tag{2.6}$$

and K is the Kobayashi-Maskawa matrix and M_u and M_d are the diagonal mass matrices for the charge $\frac{2}{3}$ and charge $-\frac{1}{3}$ quarks.

This model may be constructed naturally by imposing the symmetry $D_R \rightarrow -D_R$ and $H_2 \rightarrow -H_2$ on the Lagrangian. The Lagrangian of Eq. (2.5) is diagonal in flavor space and so avoids problems with flavor changing neutral currents. In this model, however, the $H_2^{\,+}$ will decay preferentially to the heaviest quark. In Section III, we see that in order to explain the monojets in terms of Higgs production it is necessary that the $H_2^{\,+}$ decay predominantly to $\tau^+\nu_\tau$ (or LN) and have a large coupling to $u\bar{d}$. Hence no model in which

the Higgs Yukawa couplings are proportional to the fermion mass will satisfy our purpose.

We consider next an alternate model in which both the charge $\frac{2}{3}$ and the charge $-\frac{1}{3}$ quarks get their mass from coupling to $\tilde{H}_1$. In this model $v_2 = 0$ and so the rotated basis (Eq. (2.3)) is identical to the original basis. The Yukawa couplings to the charged Higgs are then,

$$\mathcal{L} = [\bar{U}_R A D_L + \bar{U}_L B D_R] H_2^+ \tag{2.7}$$

where A and B are arbitrary 3×3 matrices which are not in general diagonal. Since the model of Eq. (2.7) does not follow from a symmetry of the Lagrangian, the parameters of A and B must be fine tuned to avoid problems with flavor changing neutral currents.

One might think that the problems encountered in the two models of Eq. (2.5) and (2.7) could be avoided by adding more SU(2) Higgs doublets. However, in any model with natural flavor conservation the Yukawa couplings will be proportional to the fermion mass.[6] We are led to study models of the type of Eq. (2.7) with all their fine tuning problems.

In Section III, we consider models in which there is more than one SU(2) Higgs doublet, $\tilde{H}_\alpha$,

$$\tilde{H}_\alpha = \begin{bmatrix} \tilde{H}_\alpha^+ \\ \tilde{H}_\alpha^0 \end{bmatrix} \qquad \alpha = 1, \ldots, N \tag{2.8}$$

Our approach is to write an effective Lagrangian in which the couplings of the Higgs bosons to the fermions are arbitrary. In a more complete

model, these couplings would be predicted and/or related to each other. We than analyze the model to find the necessary restrictions on the coupling constants in order that the model correctly describe the real world and the monojet events.

We assume that only one of the neutral Higgs bosons, H_N^0, receives a vacuum expectation value. Both the charge $\frac{2}{3}$ and charge $-\frac{1}{3}$ quarks then receive their masses from the $\tilde{H}_N$ Yukawa coupling. The physical mass eigenstates, H_α, are linear combinations of the $\tilde{H}_\alpha$ which can be determined by examining the mass terms of the Higgs potential. Since $\langle H_\alpha \rangle = 0$ for $\alpha = 1, \ldots N-1$, the coupling of the H_α to the fermions are arbitrary. The effective Lagrangian for the charged Higgs, (which is all we are interested in) is,

$$\mathcal{L}_{eff} = \sum_{\alpha=1}^{N-1} [\bar{d}_i H_\alpha^- (a_{i\alpha} + b_{i\alpha}\gamma_5) u_i + \bar{e}_i H_\alpha^- (c_{i\alpha} + d_{i\alpha}\gamma_5) \nu_i] + \text{h.c.}, \tag{2.9}$$

where $d_i(u_i)$ is a charge $-\frac{1}{3}(\frac{2}{3})$ quark, $e_i(\nu_i)$ is a charge -1 lepton (and its neutrino), and i is a generation index. We allow for the possibility of both right- and left handed neutrinos.

We turn now to a numerical study of the effects in hadronic collisions of the Lagrangian of Eq. (2.9).

III. NUMERICAL RESULTS

We attempt to explain the monojets as the resonant production of a charged Higgs boson, $H_1^{\pm}$, and its subsequent decay to LN. Here N is a neutrino and L is a heavy lepton which we will initially take to be the tau. The color and spin averaged partonic cross section for $ud \to H_1^+ \to \tau^+\nu$ is,

$$\frac{d\sigma}{dt}(u\bar{d} \to H_1^+ \to \tau^+ \nu) = \frac{(a_{11}^2 + b_{11}^2)(c_{31}^2 + d_{31}^2)}{48\pi}\left(1 - \frac{m_\tau^2}{s}\right)^2 \times \frac{1}{(s - m_H^2)^2 + \Gamma_H^2}, \qquad (3.1)$$

where s is the total parton center of mass energy, m_H is the mass of H_1, and Γ_H is the decay width of H_1,

$$\Gamma_H = \sum_i \frac{m_H}{8\pi}[3(a_{i1}^2 + b_{i1}^2) + (c_{i1}^2 + d_{i1}^2)], \qquad (3.2)$$

where we have neglected the quark and lepton masses. Using the narrow width approximation,

$$\frac{d\sigma}{dt}(u\bar{d} \to H_1^+ \to \tau^+\nu) = \frac{(a_{11}^2 + b_{11}^2)(c_{31}^2 + d_{31}^2)}{48}\left(1 - \frac{m_\tau^2}{s}\right)^2 \times \frac{\delta(s - m_H^2)}{m_H \Gamma_H} \qquad (3.3)$$

Two possible decay modes of the tau are considered: (a) $\tau \to \nu\pi$ and (b) $\tau \to$ (3 charged tracks)ν. We use the phase space Monte Carlo

SAGE to integrate the parton cross sections with the quark stucture functions and to find the final state distributions. The cross section for $u\bar{d} \to H_1^+ \to \tau^+\nu \to \pi\nu\bar{\nu}$ can be calculated analytically in the narrow width approximation and the result agrees well with that obtained using SAGE. We use the quark structure functions of Eichten et al.[7] with $\Lambda_{QCD} = .29$ GeV.

Consider the decay $\tau \to \nu\pi$. The decay width can be calculated,

$$\Gamma(\tau \to \nu\pi) = \frac{G_F^2 f_\pi^2 \cos^2\theta_c}{8\pi m_\tau}(m_\tau^2 - m_\pi^2)^2 \qquad (3.4)$$

(Note that the initial tau is polarized so that there is no spin average in Eq. (3.4)). We assume initially that there is a flavor symmetry such that $a_{i\alpha} = b_{i\alpha} = c_{i\alpha} = d_{i\alpha} = g$. Figure 1 shows the missing transverse energy spectrum from $p\bar{p} \to H_1^{\pm} \to \tau^{\pm}\nu \to \pi^{\pm}\nu\bar{\nu}$ for $m_H = 100$, 150, and 200 GeV for $g = .47$. (The quantity plotted is $[d\sigma/dp_T^m]/[B.R.(H_1 \to \tau\nu)\ B.R.(\tau \to \pi\nu)]$ so that the integral corresponds to the total production cross section of H_1). The missing energy spectum is rather flat until it reaches the kinematic cutoff at $p_T^m = m_H/2$. Clearly the couplings can be adjusted to obtain the desired rate needed to explain the monojet events. Figure 1 includes the effects of the decay matrix element of Eq. (3.4). We have checked, however, that almost identical results are obtained if we allow the tau to decay according to phase space alone.

In Figure 2, we show the missing energy spectrum for $p\bar{p} \to H_1^{\pm} \to \tau^{\pm}\nu \to$ (3 charged tracks) $\nu\bar{\nu}$. In this figure, the decay matrix element is not included. It is clear that the tau decay into four particles has more of its cross section at large p_T^m than does $\tau \to \nu\pi$. For example,

for $m_H = 150$ GeV and at $\sqrt{s} = 540$ GeV, the cross section for $p\bar{p} \to \tau\bar{\nu}$ with $p_T^m > 30$ GeV is .13 nb for $\tau \to \nu\pi$ and .22 nb for $\tau \to$ (3 charged tracks)ν.

If the tau branching ratios[8] are included, we obtain the prediction for the ratio of monojets with three charged tracks to monojets with one charged track,

$$R = \frac{\sigma(p\bar{p} \to \tau\nu \to (3 \text{ charged tracks})\nu)|_{p_T^m > 30\,\text{GeV}}}{\sigma(p\bar{p} \to \tau\nu \to \pi\nu\bar{\nu})|_{p_T^m > 30\,\text{GeV}}} \sim 3 \qquad (3.5)$$

for a 150 GeV Higgs. This is consistent with the published monojet data. (Note that the data is cut at $p_T^m = 30$ GeV).

It can be seen from Fig. 2 and Eqs. (3.2) and (3.3) that in order to obtain four events with $p_T^m > 30$ GeV having three charged tracks in an integrated luminosity of 113 nb^{-1}, we will need,

$$B.R.(H_1 \to \tau\nu) = \frac{C}{(a_{11}^2 + b_{11}^2)} \qquad (3.6)$$

where C is .1, .4, and 2.6 for m_H of 100, 150, and 200 GeV respectively at $\sqrt{s} = 540$ GeV. (Here we have used a value of 18% for B.R.($\tau \to$ (3 charged tracks)ν).) In obtaining the rates shown above we have assumed that the couplings to quarks are universal. The production rate is dominated by up and down quarks so that the couplings to the other generations are not critical.

This universality cannot be true in the case of the leptonic couplings of the H_1. If the leptonic couplings were universal then events with single electrons or single muons would occur at the same rate as those with taus. Such events would have the classic Jacobian

peak structure and would certainly have been seen if the transverse momenta of the leptons were greater than $M_w/2$; values below this could have disappeared below the signal from W decay. This would seem to eliminate universal couplings to leptons if $M_H > M_w$. We will therefore require that with $|c_{i\alpha}| = |d_{i\alpha}|$,

$$c_{11}, c_{21} < \frac{c_{31}}{\sqrt{N \cdot B.R.(\tau \to \pi\nu)}}, \qquad (3.7)$$

where N is the number of observed monojets. This criterion will ensure that there is less than one single electron or muon event.

Single electron or muon events will still arise from the decay $\tau \to e\nu\bar{\nu}$ or $\tau \to \mu\nu\bar{\nu}$ at a rate comparable to that from $\tau \to$ (3 charged tracks) + ν. The spectrum of these single leptons is shown in Figure 3. Notice that the presence of two neutrinos in the tau decay reduces the missing p_T compared with the hadronic decays. If their transverse momentum is less than $M_W/2$, the small number of events is likely to have been swamped by those from W decay,

The most stringent limits on the allowed Higgs-fermion couplings come from the tau lifetime. We consider two possible cases. In the first example, $c_{31} = d_{31}$ and so the H_1 couples only to a right handed tau neutrino. If the mass of $\nu_{\tau R}$ is greater than the tau mass then there is no contribution to the tau lifetime from Higgs exchange and this problem can be avoided. In the second example, $c_{31} = -d_{31}$ which corresponds to a left-handed neutrino. In this case, the contribution from charged Higgs exchange, H_1, is,

$$\Gamma(\tau^+ \to \nu u\bar{d}) = \frac{(a_{11}^2 + b_{11}^2)(c_{31}^2 + d_{31}^2)m_\tau^5}{512\,\pi^3 m_H^4} \qquad (3.8)$$

where we have taken $c_{11} = c_{21} = d_{11} = d_{21} = 0$. (Note that there is no interference between the contributions from charged Higgs and W boson exchange in the limit of massless up and down quarks). We will require that the contribution to the tau width from Higgs exchange be less than 1.93×10^{-13} GeV, (this is 10% of the measured value), which yields the restriction,

$$(a_{11}^2 + b_{11}^2)(c_{31}^2 + d_{31}^2) < 2.16 \times 10^{-10} \left(\frac{m_H^4}{\mathrm{GeV}}\right) \tag{3.9}$$

In order to obtain the correct number of monojets, the branching ratio $H_1 \to \tau\nu$ is contrained as in Eq. (3.6). The branching ratio is,

$$\begin{aligned} \mathrm{B.R.}(H_1 \to \tau\nu) &= \frac{c_{31}^2 + d_{31}^2}{3(a_{11}^2 + b_{11}^2) + (c_{31}^2 + d_{31}^2)} \\ &= \frac{C}{(a_{11}^2 + b_{11}^2)}, \end{aligned} \tag{3.10}$$

where C is known for a given Higgs mass and we have taken $a_{21} = a_{31} = b_{21} = b_{31} = 0$. Combining Eqs. (3.9) and (3.10) gives the restriction,

$$\frac{n}{2c}\left(\frac{m_H}{\mathrm{GeV}}\right)^4 (1-\alpha) < c_{31}^2 + d_{31}^2 < \frac{n}{2c}\left(\frac{m_H}{\mathrm{GeV}}\right)^4 (1+\alpha) \tag{3.11}$$

where $n = 2.16 \times 10^{-10}$ and,

$$\alpha = \left[1 - \frac{12C^2}{n}\left(\frac{\mathrm{GeV}^4}{m_H}\right)\right]^{1/2} \tag{3.12}$$

Using Eq. (3.12) and the numerical results for C, we see that the contribution to the tau lifetime from charged Higgs exchange is too large for all values of the Higgs mass which we have considered. We are forced to conclude that our model requires the existence of a right-handed neutrino with a mass greater than the tau mass and that the H_1 couple to $\nu_{\tau R}$ but not $\nu_{\tau L}$. An alternative solution to the τ lifetime problem is for the relevant decay to be $H_1 \to LN$ where L is a new heavy lepton of mass greater than 20 GeV. Figure 4 shows the distribution in p_T^m for $L \to$ (3 charged tracks)ν when $m_L = 30$ GeV. Of course, in this case the average charged multiplicity and invariant mass of the monojets will be larger than those resulting from τ decay. L cannot be too heavy since the decay $L \to Nq\bar{q}$ will eventually result in either very broad monojets or dijets when m_L/p_T become of order one.

The H_1 can decay back into a quark antiquark pair producing a two jet final state. In principle, the H_1 could be detected in this final state as a peak in the two jet invariant mass. In order to estimate the rates, we will assume that $c_{11} = c_{21} = 0$ and $c_{31} = d_{31} = a_{11} = b_{11} = .47$. The production of jets via $H_1^{\pm}$ is then at a rate of 2.1, .3 and, 3.7×10^{-2} nb for H masses of 100, 150, and 200 GeV, respectively at $\sqrt{s} = 540$ GeV. The continuum (QCD) production of jets is at a rate for $d\sigma/dM$ (where M is the jet pair mass and each jet has rapidity y satisfying $|y| < .85$) of 1.3, .07, and 5×10^{-2} nb/GeV at M = 100, 150, and 200 GeV. The observability of the H_1 decay to two jets depends critically on the detector resolution which can be obtained on the jet pair mass. For comparison, it has not yet proved possible to see the W boson via its hadronic decays.[9] In this case the jet rate from resonances is 3 nb, while the corresponding jet background is 5.4 nb/GeV.

One hadronic decay of the H_1 may be easier to see. The charged Higgs can decay into a $t\bar{b}$ quark pair. If the couplings of the top and bottom quarks to the charged Higgs are identical to those of the up and down quarks ($a_{31} = a_{11}$, $b_{31} = b_{11}$), then the number of top quarks, n_t, produced from H_1 decay is,

$$N_t = 3N/B.R.(\tau \rightarrow (3 \text{ charged tracks})\nu \text{ or } \pi\nu) \qquad (3.13)$$

where N is the number of monojets. The top quarks will decay semileptonically and may be included in the top search at UA1. If the semileptonic branching ratio of the top is included, we will have of order 6 events. It is not clear whether they would have been detected.

IV. CONCLUSIONS

The kinematics of the monojet events seen at UA1 are consistent with the resonant production of a heavy colorless object which then decays sequentially. A 150 GeV Higgs boson with non-standard couplings to the quarks and leptons can correctly reproduce the missing transverse energy spectrum of the monojet data. In addition, if the Higgs boson has enhanced couplings to the tau, then the low charge multiplicity and low effective mass of the monojets can also be explained.

The simplest version of the model where H decays to $\tau\ \nu_L$ is excluded by constraints from the measured value of the τ lifetime. The decay $H \rightarrow \tau\nu_R$ with $m(\nu_R) > m_\tau$ solves this problem. Alternatively if $H \rightarrow LN$ where L is a new heavy lepton , then the invariant mass and multiplicity of the monojets will be larger. New data from UA1[10] indicate that there may be more monojets. These appear to have larger multiplicity and jet mass than the earlier set.[1] In addition, a number of dijet events with two jets and missing transverse momentum have been observed. These dijets are not dominantly of the type expected by supersymmetry which predicts that the jets should be on opposite sides.[11]

Our model has two sources for dijets, the simplest being the QCD process which produces H plus jet as a final state. In a model of this type, it is also likely that $H^0 \rightarrow L^+L^-$ (or $\tau^+\tau^-$) would occur at a reasonable rate. This is a potential source of dijets; unfortunately they will be dominantly of the back to back type.

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FIGURE CAPTIONS

Fig. 1. Missing p_T spectrum for $p\bar{p} \to H_1 \to \tau\nu \to \pi\nu\bar{\nu}$. The quantity plotted on the ordinate is $d\sigma/dp_T^m/[B.R.(H_1 \to \tau\nu)\ B.R.(\tau \to \pi\nu)]$. The solid, dashed, and dot-dashed lines have m_H = 100, 150, and 200 GeV respectively and $\sqrt{s}$ = 540 GeV. The dotted, solid, and dashed lines have m_H = 100, 150, and 200 GeV respectively and $\sqrt{s}$ = 620 GeV.

Fig. 2. Missing p_T spectrum for $p\bar{p} \to H_1 \to \tau\nu \to$ (3 charged tracks)$\nu\bar{\nu}$. The quantity plotted on the ordinate is $d\sigma/dp_T^m/[B.R.(H_1 \to \tau\nu)B.R.(\tau \to$(3 charged tracks)$\nu)]$. The solid, dashed, and dot-dashed lines have m_H = 100, 150, and 200 GeV respectively and $\sqrt{s}$ = 540 GeV. The dotted, solid, and dashed lines have m_H = 100, 150, and 200 GeV respectively and $\sqrt{s}$ = 620 GeV.

Fig. 3. Missing p_T spectrum for $p\bar{p} \to H_1 \to \tau\nu \to e\bar{\nu}\nu\bar{\nu}$. The quantity plotted on the ordinate is $d\sigma/dp_T^m/[B.R.(H_1 \to \tau\nu)B.R.(\tau \to e\bar{\nu}\nu)]$. The solid, dashed, and dot-dashed lines have m_H = 100, 150, and 200 GeV respectively and $\sqrt{s}$ = 540 GeV. The dotted, solid, and dashed lines have m_H = 100, 150, and 200 GeV respectively and $\sqrt{s}$ = 620 GeV.

Fig 4. Missing p_T spectrum for $p\bar{p} \to H_1 \to LN \to$ (3 charged tracks)$\nu\bar{\nu}$. The quantity plotted on the ordinate is $d\sigma/dp_T^m/[B.R.(H_1 \to \tau\nu)B.R.(\tau \to$(3 charged tracks)$\nu)]$. The solid, dashed, and dot-dashed lines have m_H = 100, 150, and 200 GeV respectively and $\sqrt{s}$ = 540 GeV. The dotted, solid, and dashed lines have m_H = 100, 150, and 200 GeV respectively and $\sqrt{s}$ = 620 GeV.

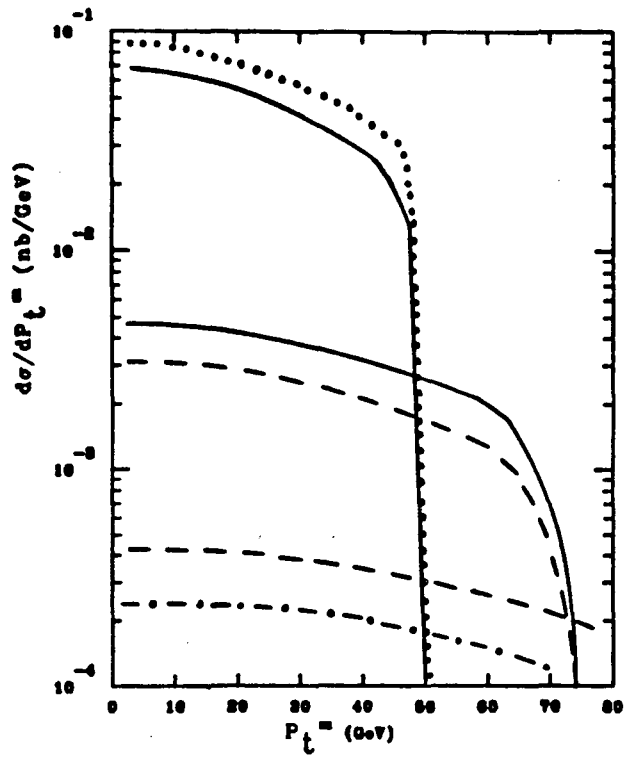

Figure 1

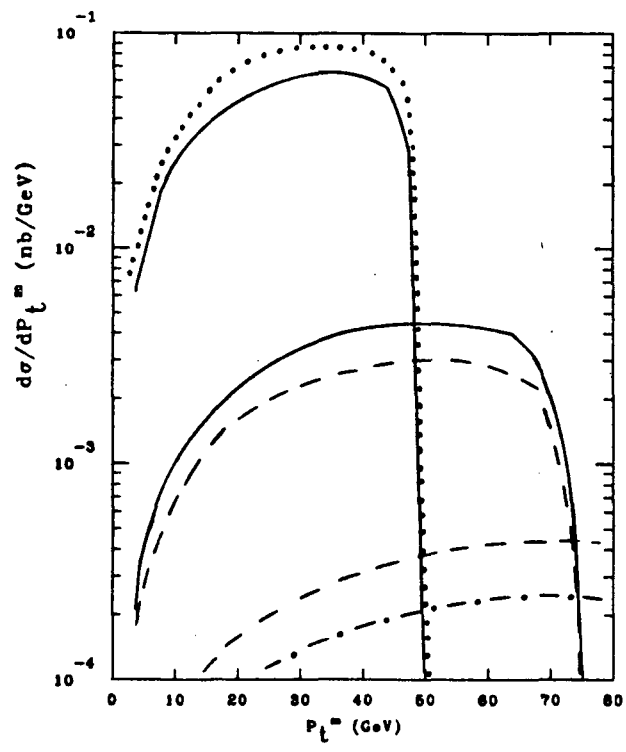

Figure 2

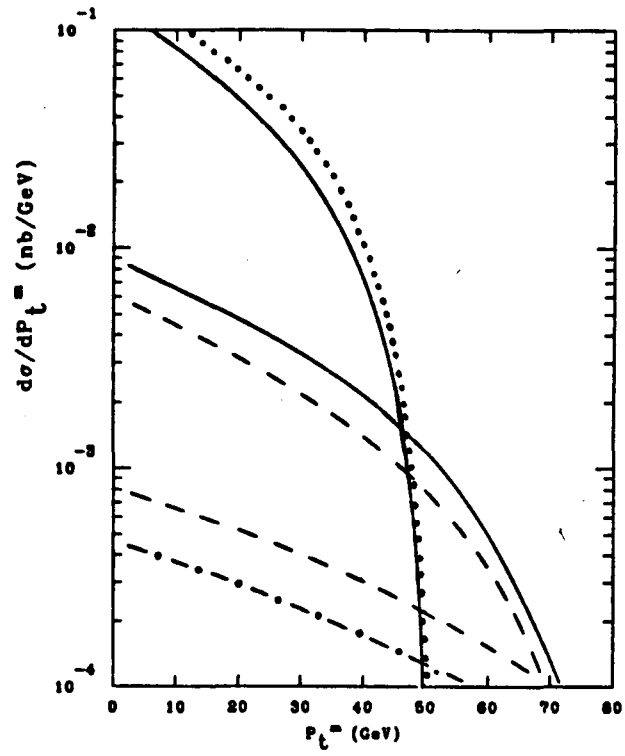

Figure 3

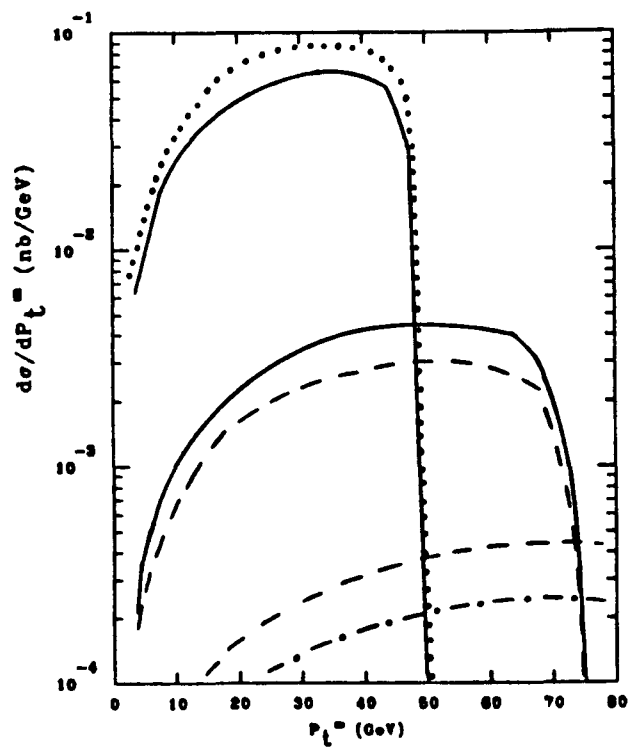

Figure 4

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