Model of New Scalar Bosons as the Source of One-Jet Events at the $\bar{p}p$ Collider

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Received 2 January 1986

Abstract. A model involving color octet and singlet scalar bosons is studied as a possible candidate to account for the one-jet events with large missing transverse energy observed at the CERN proton-antiproton collider. A distinguishing characteristic of this model, as compared to various models based on supersymmetry, is that the heavy particle being produced is assumed to be also the carrier of the missing momentum.

Some events have been observed \cite{1} in proton-antiproton collisions to consist of nothing but a single jet of particles produced at high transverse momentum. The large missing transverse energy is presumably due to one or more particles which escape the detector unobserved. These events are called monojets, and they have received a lot of attention theoretically because they may be a signal of new physics beyond the standard model. Among the many models that have been proposed, those based on supersymmetry are the most vigorously pursued \cite{2-7}. In all these explanations of the monojets, heavy color octet fermions (gluinos) or heavy color triplet scalar bosons (squarks) are produced, and the missing transverse energy is carried away by light color singlet fermions (photinos) among the decay products. In this paper, we explore a rather different situation, not necessarily based on supersymmetry, where light color octet scalar bosons are produced together with a heavy color singlet scalar boson which decays partially into other particles that escape the detector unobserved. Theoretically, the existence of light color octet scalar bosons is a natural consequence of a model of global $SU(3)$ color \cite{8}, whereas the possibility that a neutral Higgs boson will sometimes be invisible in its decay is certainly not without foundation, as has been pointed out recently \cite{9}. In any case, our model offers a rather different phenomenology from that of all others, and should at least be useful for the purpose of comparison as more experimental data are being acquired.

Let $x_a$ ($a=1,\ldots,8$) be an octet of light scalar bosons which transform as $(8,1,0)$ under $SU(3)\times SU(2)\times U(1)$ of the standard model. Then their locally gauge-invariant interaction with the vector gluons $A^a_\mu$ is given by

$$\mathcal{L}_{\text{int}} = g f_{abc} A^a_\mu (\partial_\mu x_b) x_c + \frac{1}{4} g^2 (f_{ace} f_{bcd} + f_{ade} f_{bce}) A^a_\mu A^b_\nu x_c x_d,$$  \hspace{1cm} (1)

where $g$ is the usual $SU(3)$ coupling of quantum chromodynamics, and $f_{abc}$ are the conventionally defined $SU(3)$ structure constants. Since the $x_a$'s are $SU(2)$ singlets, they do not couple to the known quarks, hence their effect on ordinary hadron dynamics is mostly negligible. In any given process of perturbative quantum chromodynamics initiated by quarks, they are produced only through their interaction with the vector gluons. Therefore, their presence is only felt through higher-order corrections which are not as severely tested by experiment. They are also electrically neutral, so they are not directly produced in $e^+ e^-$ colliders. However, if there is sufficient gluonic content in the initial states, such as in a high-energy $\bar{p}p$ collision, they can be produced rather copiously and observed as jets. At lower energies, they are also produced, perhaps with a cross section similar to that of charm production, which is again not completely understood.
Let $H$ be a heavy neutral scalar boson which is a color singlet and couples to $x_a$ according to

$$\mathcal{L}'_{\text{int}} = f v H(x_a x_a) + \ldots ,$$  

where $f$ is in principle arbitrary, but will be assumed equal to the $SU(2)$ gauge coupling for calculational purposes, and $v = (\sqrt{2} G_F)^{-1/2} \approx 246 \text{ GeV}$. There are undoubtedly more terms on the right-hand side of (2), but they will not be of importance with regard to the discussion at hand. Therefore, $H$ is identifiable as a neutral Higgs boson which decays at least into two $x_a$'s. We now must assume that it also decays into other particles which escape the detector unobserved [9]. This is certainly an ad hoc assumption, but a necessary one if we are to explore the possibility that the large missing energy in a monojet event is due entirely to a new heavy particle, whatever it may be.

In a high-energy $pp$ collision, the constituent gluons carry a large fraction of the total momentum of each projectile. In the center of mass of the colliding gluons, the cross section for producing $xxH$, with a small mass $m_x$ for the $x$ boson, is given by

$$\frac{d\sigma}{dE_3 d\theta_3 d\theta_4 d\phi} = \frac{|\mathcal{M}|^2}{(4\pi)^4} \frac{s_3 s_4 E_3 [4E(E-E_3)-m_H^2]}{(4E)^2 [2E-E_3(1-C_{34})]^3}, \quad (3)$$

where $E_3$ is the energy of one of the $x$'s, $\theta_3$ its angle with respect to the beam axis, $\theta_4$ the corresponding angle of the other $x$, $\phi$ the azimuthal angle between the two $x$'s, $E$ the energy of either initial gluon, $m_H$ the mass of the $H$ boson, $s_3 = \sin \theta_3, s_4 = \sin \theta_4, c_{34} = \cos \theta_3 \cos \theta_4 + \sin \theta_3 \sin \theta_4 \cos \phi$, and

$$|\mathcal{M}|^2 = \frac{9}{16} g^4 f^2 v^2 \left\{ \begin{array}{l}
\left[ -3 \frac{6E_3}{E} + \frac{E_3^2}{E^2} (6+c_3^2) \right] + \frac{1}{4E^2(E-E_3)^2} \\
\left[ -3 \frac{6E_4}{E} + \frac{E_4^2}{E^2} (6+c_3^2) \right] + \frac{1}{4E^2(E-3E_3)^2} \\
\left[ -6 \frac{(E_3+E_4)}{E} + \frac{2E_3E_4}{E^2} (6c_3^2 + c_3^4(1-2c_3^2)) \right] + \frac{c_\phi}{s_3 s_4 E^3(E-E_3)} \left[ (1-c_3 c_4) \left( -3 + \frac{7E_3}{E} - \frac{E_3^2}{E^2} \right) \right]
\end{array} \right\}$$

$$+ \frac{c_\phi}{s_3 s_4 E^3(E-E_3)} \left[ (1-c_3 c_4) \left( -3 + \frac{7E_4}{E} - \frac{E_4^2}{E^2} \right) \right]
+ 2 \frac{c_\phi}{E^4} \left[ c_\phi^2 + 2 \left( 1 - \frac{2}{s_3} + \frac{4}{s_3^2} \right) \frac{4(1-c_3 c_4)}{s_3^2 s_4^2} \right], \quad (4)$$

with $c_3 = \cos \theta_3, c_4 = \cos \theta_4, c_\phi = \cos \phi$, and

$$E_4 = \frac{4E(E-E_3)-m_H^2}{2[2E-E_3(1-c_{34})]} \quad (5)$$

The two $x$'s will be observed as jets if they pass the appropriate transverse energy cuts: $E_T > 25 \text{ GeV}$ for the primary jet, $E_T > 12 \text{ GeV}$ for the secondary jet; and $H$ will carry away the missing energy. With the jet $E_T$ cut, the probability that the less energetic $x$ in a given event will be counted as a jet is lowered considerably. This is in accordance with what has been observed.

To compare against the actual data, we use the parton calculation as given by (3) and (4), with the gluon distribution inside the proton taken from Duke and Owens [10] with $A = 0.2 \text{ GeV}$. We adopt the jet finding algorithm as used by Barger et al. [4], but do not include fragmentation effects. Since the missing $p_T$ in our model comes from $H$ and not $x$, fragmentation will not significantly affect that distribution. However, since $m_x$ is assumed to be small, the initial proton and antiproton themselves may contain a nonnegligible fraction of $x$'s. Hence the process $x + \text{gluon} \rightarrow x + H$ may contribute significantly to the monojet cross section, especially at large missing $p_T$. In fact, we can use the evolution equations of quantum chromodynamics to come up with a distribution of the $x$ boson inside a nucleon for a given value of $m_x$. We have calculated the splitting functions involving $x$, and they are given by

$$P_{gs} = 6 \left[ \frac{1-z}{z} \right], \quad (6)$$

$$P_{sg} = 6z(1-z), \quad (7)$$

$$P_{xs} = 6z \left[ \frac{6z}{(1-z)^4} + 6 \delta (1-z) \right], \quad (8)$$

where $g$ refers to the gluon, and $z$ is equal to the fraction of momentum carried by the particle which splits off. The splitting function $P_{xs}$ in our case differs from that of the standard model in its regularization at $z = 1$, namely

$$P_{xs} = 6 \left[ \frac{1-z}{z} + \frac{z}{(1-z)^4} + z(1-z) + \frac{11}{18} \delta (1-z) \right], \quad (9)$$

due to the presence of $P_{gs}$. The other quark and gluon splitting functions remain the same. We assume that the $x$-boson distribution is zero for $Q < Q_0 = 2m_x$, and take the quark and gluon distributions of Duke and Owens at $Q_0$ as inputs. Above $Q_0$ the running QCD coupling becomes

$$\alpha_s = \left[ 2\beta_0 \ln (\frac{Q}{A}) \right]^{-1} \quad (10)$$