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IC/78/88

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INTERNATIONAL ATOMIC ENERGY AGENCY

UNITED NATIONS EDUCATIONAL, SCIENTIFIC AND CULTURAL ORGANIZATION

1978 MIRAMARE-TRIESTE

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SIGNATURES FOR AXIAL CHROMODYNAMICS *

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MIRAMARE - TRIESTE
July 1978

* To be submitted for publication.

** Supported in part by the National Science Foundation, under Grant No.GP43662X.

ABSTRACT

Within the context of basic left-right symmetry and the hypothesis of unification of weak, electromagnetic and strong forces at a mass level $\approx 10^4$-10^6 GeV, relatively light "mass" axial gluons, confined or liberated, must be postulated. We remark that the existence of such "light" axial gluons supplementing the familiar vector octet preserves the successes of QCD, both for deep inelastic processes and charmonium physics. Through the characteristic spin-spin force, generated by their exchange, they may even help resolve some of the discrepancies between vector QCD predictions and charmonium physics. The main remark of this note is that if colour is liberated, not only vector but also axial-vector gluons are produced in high-energy e^-e^+ experiments, e.g. at PETRA and PEP, with fairly large cross-section. Distinctive decay modes of such liberated axial gluons are noted.

I. THE NEED FOR CHIRAL COLOUR

Four-component quarks possessing flavour and colour exhibit chiral flavour gauge interactions. It appears surprising to us that it has essentially been taken for granted that colour gauge interactions are vectorial only and not chiral.

A priori, there are several theoretical motivations for postulating that basic colour gauge interactions are chiral just like the flavour ones.[1] But perhaps the most important one is this; introduction of chiral colour gauges (for example,[2] in $[SU(4)]^4 = [SU(4)_A \times SU(4)_B]_{flavour} \times [SU(4)'_C \times SU(4)'_D]_{colour}$, or more generally $[SU(n)]^4$ with $n \geq 4$) together with the assumption that the quark chiral colour subgroup $SU(3)'_L \times SU(3)'_R$ is preserved as a good low-energy symmetry, permits a unification of weak, electromagnetic as well as strong interactions to become manifest at a mass scale as low as 10^4-10^6 GeV.[3] Experimentally this implies that there could be direct signatures of this unification already at Isabelle. By contrast, alternative unifying symmetries [4] (e.g. SU(5), SO(10), E_6 and E_7) devoid of chiral colour, predict ultraheavy unifying mass scales $\gtrsim 10^{15}$ GeV.

As will be elaborated later, there is nothing in present low-energy regime to exclude colour being chiral just like flavour. In fact an underlying chiral colour symmetry, by providing an octet of axial vector gluons [2], in addition to the familiar vector octet, may even help resolve some of the lingering discrepancies between predictions of vector QCD and charmonium physics.

Now chiral colour $SU(3)'_L \times SU(3)'_R$ cannot be preserved as an exact local symmetry, since quarks possess mass. We expect that it breaks at the secondary stage of SSB to vectorial colour symmetry $SU(3)'_{L+R}$ generating a massive octet of axial-vector gluons $V_A(\underset{\sim}{8})$ together with the familiar octet of vector gluons $V(\underset{\sim}{8})$. In accordance with the low mass unification hypothesis, the breaking of chiral colour must be at the mass level no bigger than the breakdown of the "electro-weak" symmetry $SU(2)_L \times U(1)$. Correspondingly, we should expect a "relatively light" octet of axial gluons with "masses" which at the one extreme may be as small as nearly 1 GeV and at the other, perhaps as high as 100 GeV. (By "mass" here, we mean the Archimedes mass [5], which the gluons would seem to possess inside hadronic bags, just as constituent (p,n) quarks inside flavour singlet baryons exhibit Archimedes' masses ~300 MeV and b-quark masses ~5 GeV.) If the gluons are liberated, their outside physical masses would, of course, be higher than their inside effective masses - see remarks later.

The purpose of this note is two-fold: (a) to make some general remarks regarding the consequences and the consistency of the hypothesis of light axial colour gluons, whether they are confined or liberated, and (b) to point out that if axial colour gluons are liberated rather than confined, they would exhibit unambiguous signatures for high-energy e^-e^+ experiments, e.g. at PETRA and PEP. They would be produced with cross-sections corresponding to electronic partial widths of order 100 keV to few MeV and can be distinguished by their rather distinctive decay modes.

II. MANIFESTATIONS OF CHIRAL COLOUR FOR CONFINED OR LIBERATED GLUONS

Within unified theories like $[SU(n)]^4$, quarks may be fractionally or[1] integrally charged. For the fractional charge case, the vector gluons are necessarily massless[6]. In this case, exact colour confinement is generally postulated; for the integer-charge case with necessarily finite mass gluons, the confinement would only be partial, with coloured quarks and gluons exhibiting the Archimedes effect[5], with different "masses" inside and outside hadronic bags.

By "inside" mass we mean the following. In a renormalizable gauge theory, with mass generation through SSB, one can define running masses $\bar{\mu}(q^2)$ as functions of momenta q^2, from renormalization group considerations. We shall define "inside" masses (m_{in}) as values of $\bar{\mu}(q^2)$ for $q^2 \approx 1/R^2$, where R is the radius of hadronic bags of the order of a Fermi. For exact confinement the outside mass (m_{out}) is infinite; for partial confinement it is finite. The precise link between m_{out} and m_{in}, involving as it does long distance, possibly non-perturbative physics, is an unsolved problem. Semi-quantitative considerations based on the picture of a partially confining pervious bag model have been offered, for example in a formula like[7] $(m_{quark})_{out} \approx \frac{1}{2\pi\alpha'\mu_V}$ where α' is the Regge trajectory slope (in GeV^{-2}) and μ_V is the inside vector gluon mass. In the sequel, for the case of partial confinement we shall, for orientation purposes, use an ansatz $m_{out} - m_{in} \approx \frac{1}{2\pi\alpha'\mu_V}$. Note that m_{in} (from SSB) will depend not only on particular (colour) multiplets but also possibly on the flavour components concerned. For example, m_{in} for (n,p) quarks is $\approx$ 300 MeV, while $(m_{in})_b$ is $\approx$ 5 GeV. With μ_V varying between 100 and 10 MeV, $m_{out}-m_{in}$ is of order $\approx$ 2-20 GeV. Since, from the low mass unification hypothesis and the SSB

implied by it, a priori, we cannot place sharper bounds on $(m_A)_{in}$ for axial gluons, except that $(m_A)_{in}$ range between nearly one and 100 GeV, we expect that $(m_A)_{out}$ would range between few to 100 GeV for the partially confined integer-charge quark case. (Since from any SSB arguments we expect $(m_A)_{in} > (m_V)_{in} \equiv \mu_V$, we expect $(m_A)_{out}$ also to be $>(m_V)_{out}$.)

But irrespective of whether axial gluons are exactly or partially confined, light axial gluons would generate an effective low-energy axial chromodynamics supplementing the familiar vector QCD. We remark that the presence of the axial gluon octet with either very light running masses $\bar{\mu}_A(\langle q\rangle^2) \lesssim$ a few hundred MeV or very heavy $\bar{\mu}_A^2(\langle q\rangle^2) \gg \langle q^2\rangle$ will not disturb the familiar successes of QCD. (Here $\langle q^2\rangle$ denotes average typical (momentum transfer)^2 for deep inelastic scaling, i.e. $\langle q^2\rangle > 1$ GeV^2.) This is because in the case of vanishing $\bar{\mu}_A(q^2)$, together with vanishing $\bar{\mu}_{quark}(q^2)$, one recovers unbroken chiral SU(3)' colour. In this case left and right quarks couple to separate octets of gluons $V_L(8)$ and $V_R(8)$ with equal coupling constants. Hence all applications of QCD remain unaltered. (With finite $\bar{\mu}_A^2(q^2 \ll \langle q^2\rangle$, departures from QCD predictions in respect of moments and structure functions would only be of order $\bar{\mu}_A^2/\langle q^2\rangle \lesssim 1\%$.)

For the case of heavy axial masses $\mu_A^2(\langle q\rangle^2) \gg \langle q^2\rangle$, on the other hand, one can effectively ignore the contributions of the axial octet altogether. Only vector chromodynamics would be relevant[3] with a coupling which is $1/\sqrt{2}$ of the chiral colour coupling. This too leaves QCD applications unaltered. Thus from the successes of QCD, we would infer that either $\mu_A^2(\langle q\rangle^2) \ll 1$ GeV^2 or that $\mu_A^2(\langle q\rangle^2) \gg \langle q^2\rangle_{max} \sim 100$ GeV^2.

Exchanges of axial gluons would however lead to a spin-spin force in the leading order in contrast to vector exchanges, where the spin-spin force is suppressed as a non-relativistic effect, particularly for low effective axial masses. The net strength of such a spin-spin force depends on $(m_A)_{in}$. With masses ~1/2-2 GeV the spin-spin force could indeed be relevant for the splitting of $^3S_1$ and $^1S_0$ $q\bar{q}$ states, e.g. (ψ(3100) and η_c(2800)) and also for radiative transitions. One must bear in mind, however, that in general, multi-axial exchanges would play an important role, as would the multi-vector exchanges. Such a spin-spin force may resolve some of the discrepancies which appear to be developing between vector QCD predictions and charmonium physics[8]. A complete calculation involving the contribution of the axial spin-spin force for the charmonium system is yet to be carried out. We urge such a calculation. This would no doubt involve, in general, at least two new parameters corresponding to the range and strength of the new spin-spin force. We stress, however, that the full potential and signature of this new force would easily be manifest (in spite of the introduction of new parameters) through a joint exploration of the spectroscopies of the charmonium as well as of the heavier bottomonium system, which should be available within the next few years.

To conclude, even for confined axial chromodynamics, one can find signatures for axial gluons [9] perhaps through departures from vector QCD predictions for deep inelastic processes, but more promisingly through an exploration of the spectroscopy of the heavy quarkonium systems. However, their most spectacular debut would be for the liberated case which we consider next.

III. LIBERATED AXIAL GLUONS

Consider the case of integer-charge quarks and liberated colour. We shall mainly be concerned with the production and decays of axial gluons in $e^+ + e^-$ collisions. The interplay between the flavour and colour gauge particles is best exhibited through the neutral mass matrix, which we now construct.

Gauge mass matrix: To be specific consider spontaneous breakdown of $[SU(4)]^4 = [SU(4)_A \times SU(4)_B]_{flavour} \times [SU(4)'_C \quad SU(4)'_D]_{colour}$. Here A,B denote left and right flavour and C,D left and right colour in the space of basic fermions. The primary breaking through VEV of 15-plets of each SU(4) reduces the symmetry $[SU(4)]^4$ to [2),3),10)]

$$\mathcal{G} = \left[SU(2)_L \times SU(2)_R\right] \times \left[SU(3)'_L \times SU(3)'_R\right] \times \left[U(1)'_L \times U(1)'_R\right] ,$$

where $SU(2)_{L,R}$ are the left and right GIM flavour subgroups and $U(1)'_{L,R}$ denote the 15^{th} generators of $SU(4)'_{L,R}$ colour. The gauge particles of $SU(2)_{L,R}$, $SU(3)'_{L,R}$ and $U(1)'_{L,R}$ are denoted in the notation [11] of Ref.2 by $(W^{\pm},W^3)_{L,R}$, $V(\underset{\sim}{8})_{L,R}$ and $S^0_{L,R}$, respectively, while their respective coupling constants are denoted by $g_{L,R} \equiv g$, $f^S_{L,R} = f_s$ and $(f_{15})_{L,R} = f_{15}$. (We are assuming that the uniconstant symmetry $[SU(4)]^4$ breaks in such a manner that left-right symmetry is preserved at the primary stage of SSB, [12] thus the radiative differences between g_L and g_R and $(f_{15})_L$ and $(f_{15})_R$ are small.)

The secondary stage of SSB is induced through fundamental multiplets. Only four such multiplets (related by discrete symmetry [12]) are needed: $A = (4,\bar{4},1,1)$, $B = (1,4,1,\bar{4})$, $C = (4,1,\bar{4},1)$, $H = (1,1,4,\bar{4})$. Assume the pattern [13] of VEV: $\langle A\rangle = (a_1,a_1,a_1,a_4)$, $B = (0,0,0,b_4)$, $\langle C\rangle = (c_1,c_1,c_1,c_4)$, $\langle H\rangle = (h_1,h_1,h_1,h_4)$, where the entries denote VEV of diagonal elements. Consistent with the low mass unification hypothesis, assume the hierarchy $b_4 >> (c_4,a_4,h_4) >> c_1$, $h_1 < h_4$ and $h_4^2 \sim (c_4^2 + a_4^2)$. The scale of these parameters is fixed by the combination [14] $c^2 + a^2 = (4m^2_{W^\pm_L}/g^2) = (\sqrt{2}G_F)^{-1}$, where $c^2 \equiv 3c_1^2 + c_4^2$ and $a^2 \equiv 3a_1^2 + a_4^2$

Note that b_4 being the largest makes the right-handed charged gauge particles $W^{\pm}_R$ and a neutral gauge particle B_R (see below) superheavy leaving an abelian $\tilde{U}(1)_R$ invariant

$$SU(2)_R \times U(1)'_R \xrightarrow{b_4} \tilde{U}(1)_R$$

$\tilde{U}(1)_R$ couples to the appropriate diagonal sum of W^3_R and S^0_R (see below). The parameter $h^2 \equiv 3h_1^2 + h_4^2$ breaks $U(1)'_L \times U(1)'_R$ to $U(1)'_{L+R}$ giving mass to the colour-singlet axial combination $(S^0_L - S^0_R)/\sqrt{2}$. The parameter h_1 breaks chiral quark colour $SU(3)'_L \times SU(3)_R$ to vector colour $SU(3)'_{L+R}$, and thereby gives mass to the octet of axial colour gluons $V_A(\underset{\sim}{8})$; while the parameter c_1 breaks left flavour $SU(2)_L$ and left colour $SU(3)_L$ (and therefore vector colour $SU(3)'_{L+R}$) giving mass to the octet of vector gluons $V(\underset{\sim}{8})$ (c_1 being the smallest of all VEV's makes the vector octet $V(\underset{\sim}{8})$ the lightest apart from the photon). The parameter c_4 breaks $SU(2)_L \times U(1)_L$ to the diagonal abelian sum $\tilde{U}(1)_L$ (analogous to (5)), while a_4 breaks flavour $SU(2)_L \times SU(2)_R$ to $SU(2)_{L+R}$. A summary of these symmetry breakings is given below:

$$U'(1)_L \times \tilde{U}(1)_R \xrightarrow{(h_4,h_1)} U(1)$$

$$SU(3)'_L \times SU(3)'_R \xrightarrow{h_1} SU(3)'_{L+R}$$

$$SU(2)_L \times U(1) \quad \times SU(3)'_{L+R} \xrightarrow{(c_4,a_4),c_1} U(1)_{EM} .$$

Here U(1) is the familiar abelian generator of $SU(2)_L \times U(1)$.

$U(1)_{EM}$ contains flavour and SU(3)' colour generators appropriate to the case of integer-charge quarks. (Note that for the case of fractionally charged quarks the above pattern of SSB applies with the exception that $c_1 = 0$.) We now present below the consequences of this breaking pattern on the neutral gauge mass matrix [15].

Neutral eigenstates: There are six relevant neutral gauge particles $W^3_{L,R}$, $S^0_{L,R}$ and $U^0_{L,R}$, which mix with each other $\left(U^0_{L,R} \equiv \frac{(\sqrt{3}\ V_3 + V_8)_{L,R}}{2}\right)$; they denote the canonical chiral gluon fields, which enter into the photon. We know the exact composition of the massless photon as well as the composition of one heavy gauge particle Z_C . They are:

$$A = \cos\phi\ (A_L + A_R)/\sqrt{2} + \sin\phi\ (U_L + U_R)/\sqrt{2}\ \ (m_A = 0)$$

$$Z_C \approx B_R \quad ; \quad m^2_{Z_C} \approx (g^2/4)\ (3r+1)\ b_4^2 \quad , \tag{1}$$

where $\tan\phi = (4r/(3r+1))^{1/2}\ (g/f_s)$, $A_{L,R} \equiv (3r+1)^{-1/2}\ (\sqrt{3r}\, W^3 - S^0)_{L,R}$ and $B_{L,R} \equiv (3r+1)^{-1/2}\ (W^3 + \sqrt{3r}\, S^0)_{L,R}$. The ratio $r \equiv f^2_{15}/2g^2$ is related to the weak angle by $\sin^2\theta_W = 3r/(6r+2)$. Extracting out the two eigenstates — photon and Z_C – we are left with a 4×4 symmetric matrix, which in units of $(g^2/4)$ is given by:

$\frac{A_L - A_R}{\sqrt{2}}$	B_L	$\bar{U}_V$	U_A
$\left(\frac{6r}{3r+1}\right)\ (h^2+a^2)$	$\sqrt{\frac{3r}{2}}\Big/(3r+1)\ (2a^2)$	$\frac{-2\sqrt{2r}}{\sqrt{3r+1}}\ (f_s/f_{15})c_1^2$	$\frac{-2\sqrt{2r}}{\sqrt{3r+1}}\ c_1^2$
	$(3r+1)^{-1}\ [9r^2h^2+a^2 + (3r+1)^2\ c_4^2]$	$-4\left[\frac{r}{3(3r+1)}\right]^{\frac{1}{2}}\left(\frac{f_s}{f_{15}}\right)c_1^2$	$-4\left[\frac{r}{3(3r+1)}\right]^{\frac{1}{2}}\left(\frac{f_s}{f_{15}}\right)\ c_1^2$
		$2r\left(\frac{f_s}{f_{15}}\right)^2\ c_1^2$	$2r(f_s/f_{15})^2\ c_1^2$
			$2r(f_s/f_{15})^2(4h_1^2+c_1^2)$

(2)

In the above we have not exhibited correction terms of order (g^2/f_s^2), (f^2_{15}/f^2_s) and $c_1^2(h^2+a^2)$ compared with unity.

Here $\bar{U}_V \equiv \cos\phi\ (U_L + U_R)/\sqrt{2} - \sin\phi\ (A_L + A_R)/\sqrt{2}$ and $U_{V,A} \equiv (U^0_L \mp U^0_R)/\sqrt{2}$. Since (c^2+a^2) is determined by the mass of W^+_L , there are three parameters which determine the complexion of the eigenstates of (2). These are: $\xi_1 \equiv c_1^2/(c^2+a^2)$, $\xi_2 \equiv h_1^2/(c^2+a^2)$ and $\xi_3 \equiv (h^2+a^2)/(c^2+a^2)$, where $h^2 \equiv (3h_1^2 + h_4^2)$. The general constraints on these ratios, in accordance with the requirements of low vector gluon mass and low mass unification hypothesis, are: $\xi_1 \ll 1$, $\xi_2 < 1$ and $1 \leqslant \xi_3 < \infty$.

To exhibit some of the main physical consequences which arise due to the presence of axial gluons, we shall assume (for illustration only) the following hierarchy consistent with the constraint mentioned above:

$$\xi_1 \ll \xi_2 \ll \xi_3, (\xi_3 \sim 1) \quad . \tag{3}$$

This amounts to assuming that $m^2_{U_V} \ll m^2_{U_A} \ll m^2_{Z_A} \ll m^2_{Z_B}$, where Z_A and Z_B are the two "weak" neutral gauge eigenstates to emerge from (2) . [Note for guidance that such a hierarchy holds for masses such as $m_{U_V} \sim 10$ GeV, $m_{U_A} \sim 30$ GeV, $m_{Z_A} \sim 70$ GeV, $m_{Z_B} \gtrsim 140$ GeV, though we must stress that in general m_{U_A} may be comparable to $m_{U_V} \approx$ few to 20 GeV on the one hand and to $m_{Z_A} \approx 60$ to 85 GeV on the other .] Subject to the above illustrative hierarchy, the eigenstates of the mass matrix are approximately given by:

$$\tilde{U}_V \approx \bar{U}_V + \mathcal{O}(\epsilon)\ U_A + \mathcal{O}(\epsilon')\ (A_L - A_R)/\sqrt{2} + \mathcal{O}(\epsilon'')\ B_L$$

$$\tilde{U}_A \approx (\cos\alpha)\ U_A + (\sin\alpha)\ (A_L - A_R)/\sqrt{2} + \mathcal{O}(\epsilon'')\ B_L + \mathcal{O}(\epsilon)\ \bar{U}_V$$

$$\tilde{Z}_A \approx \cos\beta\ \{(-\sin\alpha)\ U_A + (\cos\alpha)\ (A_L - A_R)/\sqrt{2}\ \} + (\sin\beta)\ B_L + \mathcal{O}(\epsilon'')\ U_A$$

$$\tilde{Z}_B \approx -\sin\beta\ \{(-\sin\alpha)\ U_A + (\cos\alpha)\ (A_L - A_R)/\sqrt{2}\ \} + (\cos\beta)\ B_L + \mathcal{O}(\epsilon'')\ U_A \ .$$

(4)

The masses are: $m^2_{U_V} \approx f_s^2 c_1^2/4$, $m^2_{U_A} \approx f_s^2 h_1^2$, for small [16)]
$\xi_3 = (h^2+a^2)/(c^2+a^2) \lesssim 1/2$, $m^2_{Z_A} \approx (4\sin^2\theta_W)(m^2_{W_L^\pm}\xi_3)$, $m^2_{Z_B} \approx m^2_{W_L^\pm}(3r+1)$ $(1+\mathcal{O}(\xi_3))$. For $\xi_3 >> 1$, $m_{Z_A} \to m_{Z^0} = (m_{W_L^\pm}/\cos\theta_W)$ with Z_B becoming much heavier than $W_L^\pm$.

The parameters ϵ, ϵ' and ϵ'' are small quantities given by [17)]:
$\epsilon \equiv (m_{U_V}/m_{U_A})^2$, $\epsilon' \equiv (g/f_s)(m_{U_V}/m_{Z_A})^2$, $\epsilon'' \equiv (g/f_s)(m_{U_V}/m_{Z_B})^2$. Subject to the hierarchy (8), $\epsilon' \sim \epsilon'' << \epsilon << 1$. The coefficients of the small components of the eigenstates are defined such that $|\mathcal{O}(\epsilon)| \approx |\epsilon|$ etc. The angles $\sin\alpha$ and $\sin\beta$ are given by:

$$\sin\alpha \sim (g/f_s)(m^2_{U_V}/m_{Z_A})^2 , \quad \sin\beta \sim a^2/(h^2+a^2) \quad . \qquad (5)$$

We are now in a position to observe four crucial features of this complex of eigenstates.

(A) Parity-violating component in $\tilde{U}_V$: The composition of the dominantly vector gluon $\tilde{U}_V$ is the same as in Ref.1, except for one notable change. It has picked up an axial component U_A of order $\epsilon = (m_{U_V}/m_{U_A})^2$. This is, of course, a parity-violating mixture, which could show itself in the hadron-hadron system, e.g. in the nuclear force. Let us estimate the effect. There are two types of contributions.

i) First there is the tree contribution with a single $\tilde{U}_V$ exchange between two nucleons. The contribution T' of this exchange to parity-violating amplitude is given by $T'(\text{tree}) = (f_s^2/m^2_{U_V})(\epsilon\delta^2)$, where δ is the colour octet component in the dominantly colour singlet nucleon ($\delta < 1/10$). Thus $T'(\text{tree}) \lesssim 10^{-5}\ \text{GeV}^{-2}$ for $m_{U_V} \sim 10$ GeV and $\epsilon = (m^2_{U_V}/m^2_{U_A}) \lesssim (1/10)$. We see that this amplitude is suppressed in large part due to the colour singlet nature of the nucleon.

ii) Second, there is the loop contribution with double $\tilde{U}_V$ exchange. This is a convergent loop. Even if we neglect form factors in nucleon-gluon vertex, the contribution to the parity-violating amplitude is $T'(\text{loop}) \sim (\alpha_s^2/m^2_{U_V})(m_{U_V}/m_{U_A})^2$. Given the mass ranges discussed above together with

$\alpha_s \sim 0.2$, we obtain $T'(\text{loop}) < 4 \times 10^{-5}\ \text{GeV}^{-2}$. If we use form factors, $T'(\text{loop})$ will be suppressed by another power of $(m_{U_V}/1\ \text{GeV})^{-2}$. Thus, in spite of the sizeable mixing between U_V and U_A, its parity-violating effects for colour singlet hadrons appear to be suppressed numerically to the level of G_{Fermi}.

Parenthetically, we observe the following feature. Quark-antiquark (or quark-quark) scattering can proceed, of course, through single $\tilde{U}_V$ exchange. The force thus generated will involve a parity-violating component due to U_V-U_A mixing. In case the $q\bar{q}$ or (qq) system is <u>inside</u> a normal hadron, the parity-violating amplitude thus arising would be given by "$[f_s^2/(\bar{m}_{U_V})^2]$" $(\bar{m}_{U_V}/\bar{m}_{U_A})^2$, where the bar denotes "inside" Archimedes' masses. Relative to the parity-conserving amplitude $\sim$ "$[f_s^2/\bar{m}_{U_V})^2]$", the parity-violating amplitude at the quark level is suppressed by the factor $(\bar{m}_{U_V}/\bar{m}_{U_A})^2 \lesssim 10^{-4}$ for $\bar{m}_{U_V} \sim 10$ MeV and $\bar{m}_{U_A} \gtrsim 1$ GeV.

It is not clear how to detect such a possibly large ($\sim 10^{-4}$) parity violation for the qq and $q\bar{q}$ systems when they are constituents inside a hadron, except possibly through the effects of the corresponding potentials on energy levels of nuclei. At any rate, if $\bar{m}_{U_A}$ were larger than 3 GeV, the effect would be rather small even at the quark level.

(B) Leptonic coupling of the axial gluon $\tilde{U}_A$: The state $\tilde{U}_A$ is dominantly the axial colour gluon U_A. However, it necessarily couples to leptonic $\bar{e}e$ and $\bar{\mu}\mu$ currents through the small components $(A_L-A_R)/\sqrt{2}$, B_L as well as the more important component $\tilde{U}_V$. (Note that $\tilde{U}_V$ contains the vector flavour component $(A_L+A_R)/\sqrt{2}$ of an amount $\sin\phi \approx (g/f_s)$.) Its coupling to leptons is given by

$$\mathcal{L}(\tilde{U}_A \to \bar{\ell}\ell) = \tilde{U}_{A\mu}\left[g_V^A\,\bar{e}\gamma_\mu e + g_A^A\,\bar{e}\gamma_\mu\gamma_5 e\right] + (e \to \mu) \quad , \qquad (6)$$

where

$$|g_V^a| \approx |\epsilon \sin\phi(e)| \approx 2|e^2/f_s|\ (m_{U_V}/m_{U_A})^2 \qquad (7a)$$

and

$$|g_A^A| \sim |(\sin\alpha)e| \approx 2|e^2/f_s|\ (m_{U_V}/m_{Z_A})^2 \quad . \qquad (7b)$$

The factor 2 arises by substituting [17] $g \approx 2e$. If $m_{Z_A}^2 >> m_{U_A}^2$, we see that the vector coupling will dominate over the axial-vector coupling. Note that the small but finite leptonic coupling of $\tilde{U}_A$ is a direct consequence of the gauge nature of the basic theory as well as of the photon being generated through SSB so as to carry flavour as well as colour components. The coupling (6) would lead to a leptonic partial width given by;

$$\Gamma(\tilde{U}_A \to e^+e^-) \approx (m_{U_V}/m_{U_A})^4 \; (m_{U_A}/1 \text{ GeV}) \; (75 \text{ keV}/(f_s^2/4\pi)) \qquad (8)$$

If we take $(m_{U_V}/m_{U_A})^2 \approx (1/2 \text{ to } 1/10)$, $m_{U_A} \approx 30$ GeV and $(f_s^2/4\pi) \approx 0.2$ at m_{U_A}, we get $\Gamma(\tilde{U}_A \to e^+e^-) \approx (100$ keV to 3 MeV), which relatively speaking, is rather large.

(C) Decay and production of the axial gluon $\tilde{U}_A$: The allowed decay modes of $\tilde{U}_A$ with the corresponding orders of magnitude of the amplitudes are listed below:

Amplitude

$$\left.\begin{array}{l}\tilde{U}_A \to e^-e^+,\ \mu^-\mu^+ \\ \quad \to \text{normal hadrons}\end{array}\right\} \approx 2(e^2/f_s)\,(m_{U_V}/m_{U_A})^2 \qquad (9a)$$

$$\tilde{U}_A \to \tilde{U}_V + \gamma \qquad \mathcal{O}(e) \qquad (9b)$$
$$\hookrightarrow e^-e^+,\ \mu^-\mu^+,\ \text{hadrons} + \gamma \ (\text{see Ref. 5})$$

$$\tilde{U}_A \to \tilde{U}_V + (\omega, \phi, K\bar{K}, 3\pi \text{ or } \omega + \eta) \qquad (\mathcal{O}(\text{strong})) \qquad (9c)$$
$$\hookrightarrow e^-e^+,\ \mu^-\mu^+,\ \text{hadrons} + \gamma,\ \text{etc. (Ref. 5)}.$$

$$\tilde{U}_A \to q + \bar{q} \qquad \mathcal{O}(\text{strong}) \ (\text{if } m_{U_A} > 2\, m_q) \qquad (9d)$$
$$\hookrightarrow \nu + \text{mesons (see Ref.18)}.$$

Here we have assumed that the vector gluon $\tilde{U}_V$ is lighter than the axial gluon $\tilde{U}_A$ in accordance with our previous discussions. We have in listing the above decay modes taken into account the facts that (i) $I = 0$ for $\tilde{U}_V$ and $\tilde{U}_A$, (ii) $C = -1$ and $+1$ for $\tilde{U}_V$ and $\tilde{U}_A$, (iii) $J^P = 1^-$ and 1^+ for $\tilde{U}_V$ and $\tilde{U}_A$, (iv) they are both colour octets but flavour singlets and (v) both ω and ϕ have SU(3) singlet components. In addition to the decay modes listed above, if the axial gluon $\tilde{U}_A$ has a mass more than twice that of the charged vector gluons $V^\pm$, it will decay through its $\tilde{U}_V$ component [19] into a pair of V^+V^-. The corresponding amplitude is of order $(m_{U_V}/m_{U_A})^2$, which is thus suppressed relative to (9c) and (9d). We therefore expect masses permitting the decay modes (9c) and (9d) or, at the very least (9b), to be the dominant decay modes of the axial gluon $\tilde{U}_A$. As long as a modest Q value ($\gtrsim$ few hundred MeV) is available for the decay modes (9c), we expect that the width of $\tilde{U}_A$ would be (conservatively) greater than 10 MeV, but perhaps as large as 100 to a few hundred MeV. Taking $\Gamma(\tilde{U}_A \to e^-e^+) \approx 100$ keV to 1 MeV, we thus expect leptonic branching ratio of $\tilde{U}_A$ to lie between nearly 10^{-3} and 10^{-1}.

Note that a major advantage of producing the axial gluon $\tilde{U}_A$ is that it in turn becomes a dominant source of the vector gluon $\tilde{U}_V$ and possibly also of quarks and antiquarks. (Decay modes of $\tilde{U}_V$ and of quarks are listed in Refs.5 and 18, respectively.)

Referring to the production of $\tilde{U}_A$, it should eventually be produced, of course, in high-energy hadronic collisions in association with other colour octet objects. But associated production cross-sections of such heavy objects would presumably be much below the level of Υ production. Fortunately, due to the leptonic coupling of $\tilde{U}_A$ of order $2(e^2/f_s)\,(m_{U_V}/m_{U_A})^2$, which yields a leptonic partial width $\Gamma(\tilde{U}_A \to e^-e^+) \approx 100$ keV to few MeV (for $m_{U_A} \approx 30$ GeV, see Eq.(8)), we expect that $\tilde{U}_A$ would be produced with a fairly decent cross-section as a resonant particle in high-energy e^-e^+ collision. In other words, if m_{U_A} lies between 10 and 40 GeV, PETRA and PEP would be ideally suited to produce the liberated axial gluon $\tilde{U}_A$ and thereby the vector gluon $\tilde{U}_V$ as well as possibly liberated quark-antiquark pairs.

What is striking is that it is very difficult for any other particle we can think of, to mimick the pattern of decay modes listed above. A weak gauge Z^0-like particle would in general decay to leptons and hadrons with similar partial widths and, most important, neither a Z^0-like particle [20] nor a

$q\tilde{q}$ resonance involving either old or new flavours would have a special affinity to produce the vector gluon $\tilde{U}_V$ via its decay modes. Thus once the axial gluon $\tilde{U}_A$ is produced in e^-e^+ collision (and it would have to be at the appropriate energy if the idea of liberation is correct), it would be difficult to identify this particle with any object other than the axial gluon $\tilde{U}_A$.

In summary, the presence of axial chromodynamics with either confined or liberated colour is important for low mass unification. Signatures for such axial chromodynamics would be apparent most promisingly from exploration of the spectroscopies of both the charmonium and the bottomonium systems. Electron-positron accelerators, e.g. PETRA and PEP or conceivably successors thereof, give the best promise for producing the liberated axial gluon $\tilde{U}_A$. Discovery of liberated axial gluons would make chiral colour and simultaneously liberated colour compulsive concepts in particle physics.

We thank Victor Elias and Subhash Rajpoot for helpful discussions.

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9) In this note we have not considered new features brought about (a) in Zweig-like selection rules and (b) in jet phenomena, where charged axial gluons with $C = +1$ would mark their presence.

10) For different complexions in the breakdown of $SU(2)_L \times SU(2)_R \times U(1)'_L \times U(1)'_R$ without chiral breaking, see Q. Shafi and Chr. Wetterich, Phys. Letters 73B, 65 (1978); V. Elias, J.C. Pati and Abdus Salam, Phys. Letters 73B, 451(1978); J.C. Pati and S. Rajpoot, ICTP, Trieste, preprint IC/78/1 and Q. Shafi and Chr. Wetterich, University of Freiburg preprint (June 1978).

11) For canonical couplings of the gauge fields see the second papers of Ref.1 and Ref.10.

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13) A natural realization of the first stage of this hierarchy, i.e. $b_4 \gg (c_i, a_i, h_i)$ can be shown to arise through a minimum of the potential (Ref.12). Whether the subsequent stage, e.g. $c_4^2 + a_4^2 \gg c_1^2 \neq 0$ can be realized through the radiative or non-perturbative origin remains to be examined. Notwithstanding the delicacy of the Higgs system, we shall proceed tentatively assuming such a hierarchy for phenomenological purposes.

14) In general, W_L's may get mass almost entirely from a^2 rather than c^2.

15) In the charged sector, charged vector and axial vectors will mix with $W_L^{\pm}$ (see Ref.1). They would also mix with each other to an extent similar to U_V-U_A mixing as discussed in the text with similar consequences.

16) Expressions for the masses of $\tilde{Z}_A$ and $\tilde{Z}_B$ for general ε_3 are given in the paper by J.C. Pati and S. Rajpoot, Ref.10.

17) If we had taken for generality $\langle B \rangle = (b_1, b_1, b_1, b_4)$, the parity violating parameters ϵ and ϵ' given in the text would be multiplied by a factor $(c_1^2 - b_1^2)/(c_1^2 + b_1^2)$. In defining these small parameters, we regard $\mathcal{O}(g/f_s) = \mathcal{O}(f_{15}/f_s)$. Note that for $\sin^2\theta_W = 1/4$ and $(f_s^2/4\pi) \approx 0.2$ (at the physical vector gluon mass) we have $g^2 = (3/2) f_{15}^2 \approx 4e^2$.

18) See J.C. Pati, Abdus Salam and S. Sakakibara, Phys. Rev. Letters 36, 1229 (1976), and J.C. Pati and Abdus Salam, ICTP, Trieste, preprint IC/78/54 (to appear in Nucl. Phys.) for a listing of quark decay modes.

19) It cannot decay through its dominant U_A component into V^+V^- because of charge conjugation plus parity invariance.

20) A weak Z^0 in general would also lead to parity-conserving forward-backward asymmetry in $e^-e^+ \to \mu^-\mu^+$ through interference between vector and axial-vector couplings, which is prominent ($\gtrsim 15\%$) over a wide range of centre-of-mass energy $E_{CM} \sim (m_{Z^0}/2)$. It is easy to see, by using Eq.(6), that $\tilde{U}_A$ exchange by contrast leads to F-B asymmetry, which is much too low ($\ll 1\%$) for $|E_{CM} - m_{U_A}| \gtrsim 10\,\Gamma(U_A) \lesssim 1$ GeV, and that even for $E_{CM} \approx m_{U_A} \pm \Gamma(U_A)$ it is at most of order 1%.

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