Disoriented Chiral Condensates:
A White Paper for the Full Acceptance Detector

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ABSTRACT

Theoretical speculations and experimental data suggesting the possibility of observing disoriented chiral condensates at a Full Acceptance Detector are reviewed.
1. Introduction

These notes are intended as an introduction to the theory and phenomenology of disoriented chiral condensates. Briefly speaking, we are concerned with the possibility that coherent states of pions might be produced in hadron-hadron collisions in existing or planned particle accelerators. Such phenomena would be signaled by events with anomalously large (or small) numbers of $\pi^0$'s in comparison with the number of charged pions. One purpose of these notes is to attempt to estimate the cross section for production of such events, the expected distribution of the neutral pion frequency, and the expected momentum distribution of such events. In the process of developing these estimates we will review theoretical grounds for believing that such events may occur. We shall also attempt to confront these theoretical ideas with available data, concentrating on the phenomenology of Centauro-type events. We hope that before reaching the end of these notes the reader agrees with us that this is an area which is fascinating both theoretically and experimentally, and that the interest is such as to warrant mounting additional efforts to observe the associated phenomena.

Before beginning, we note that most of the basic ideas contained here will be found in the various papers Bjorken has written in the past year in support of a Full-Acceptance Detector $^{1-4}$. This present work is mainly designed to try to develop the ideas a bit further, and attempt to confront the ideas with the occurrence of Centauro events and related phenomena.

First, what is a “disoriented chiral condensate?” This requires a brief digression into the physics of chiral symmetry.

Consider QCD with an isospin doublet of massless quarks

$$\psi = \begin{pmatrix} u \\ d \end{pmatrix}.$$  

The fermion terms in the Lagrangian are invariant under both global isospin transformations

$$\delta \psi = i \epsilon \cdot T \psi$$

and under global chiral transformations

$$\delta \psi = i \gamma_5 \epsilon \cdot T \psi.$$

It is widely assumed that the chiral symmetry is spontaneously broken in the QCD ground state, leading to the interpretation of the pion as a massless Goldstone
boson, and generating a mass for the quarks in the process. (In the real world, there are, apparently, small quark masses explicitly breaking chiral symmetry in the QCD Lagrangian, leading to the small but non-vanishing pion mass.) A signal of the spontaneous breaking of chiral symmetry is the non-vanishing of the vacuum condensate $\langle \bar{\psi} \psi \rangle$.

The physics which we are studying here corresponds to the possibility that in a macroscopic (but localized) region of spacetime, the vacuum condensate may be disoriented from its external or ambient value, so that the condensate is associated with the isovector, rather than the isoscalar, degrees of freedom.

Our first theoretical problem is one of a choice of language. While we could, in principle, work in the full theory of QCD, this appears to be much too complicated. Instead, since at least the late stages of the evolution of the chiral condensates are expected to involve characteristic momenta which are small compared to $\Lambda_{QCD}$, we find it more convenient to discuss the phenomena in terms of a low energy effective theory: the sigma model. At various times in our discussion we will be faced with the question of which variation of the sigma model is appropriate. For example, if we want to consider intermediate states in which a region of spacetime is locally in a chirally symmetric phase, then we will adopt the linear sigma model. We will also from time to time want to include quarks in our discussion, in which case we shall resort to the chiral quark model, in either its linear or non-linear incarnation, as appropriate. We should note at the outset that the use of the non-linear sigma model to calculate low energy processes is well established. Incorporating quarks in this version of the model (the chiral quark model) also has some successes, but is beset by questions relating to the potential double-counting of the pion. In discussions of the chiral phase transition, the utility of the linear sigma model has again been recently urged; the validity of the linear, chiral quark model has also been argued, although this appears to be rather controversial. Finally, on occasion we will include explicit chiral symmetry breaking in the form of a non-vanishing pion mass.

The upshot of all of this is that we appear to have few quantitative theoretical tools at our disposal, particularly in addressing questions related to the production of disoriented chiral condensates. Nevertheless, we do have more or less reliable qualitative and semi-quantitative tools for describing disoriented chiral condensates, particularly in the late stages of their evolution.

At this point, it may be useful to introduce a simple model for the production of disoriented chiral condensates. In this picture, we imagine that in some fraction of the collisions between ‘constituent quarks’ at sufficiently high energies, a significant fraction of the incident momentum is thermalized, resulting in small regions
of extremely hot hadronic matter, presumably hot enough to be in a chirally symmetric phase. We then expect the hadronic debris to expand outwards (transverse flow is important!) at a speed approaching that of light leaving the interior to cool rapidly. As the interior thus drops below the critical temperature, it will undergo a phase transition to a broken phase. What is crucial for our purposes is that the order parameter characterizing this phase need not have the same value that it does in the rest of the universe since the interior is protected from the exterior by a hot shell of rapidly expanding hadronic matter that is presumably in a chirally symmetric phase. Eventually, the shell cools to the hadronization temperature, and the interior re-establishes contact with the rest of the universe. At this point, the interior will realign with the exterior through the coherent radiation of non-relativistic pions. It is amusing to note the similarities between this picture and the ‘Baked Alaska’ model for nucleation of phase transitions in superfluid He-3\textsuperscript{11}.

We consider the preceding scenario to be suspect in many regards. It is far from clear that such a classical picture is appropriate, particularly in the early phases. Indeed, it is not clear that the process of creating disoriented chiral condensates need even proceed via an intermediate state dominated by a localized symmetric vacuum state. Nevertheless we believe that it is a useful starting point for discussing some of the anticipated phenomenology of disoriented chiral condensates.

In this context, a working definition of an event involving a disoriented chiral condensate is:

(a) At a late stage in the production process, just before the relaxation of the inside to the outside vacuum, there is (in a suitable frame) a roughly spherical volume of disoriented chiral condensate enveloped by a uniformly packed shell of hadrons that are moving outward at the speed of light.

(b) The localized chiral condensate, (the “inside vacuum”) then relaxes to the outside vacuum accompanied by the radiation of Goldstone bosons, viz., primarily pions.

The relevant physics then involves three fundamental issues:

(i) Is there a threshold realizable at existing or currently planned hadron colliders for producing a disoriented chiral condensate and what are testable models of possible production mechanisms in pp collisions?

(ii) What are the orders of magnitude of the cross sections for the occurrence of processes in which the disorientation and subsequent relaxation of the chiral condensate plays a central role?

(iii) What are the expected experimental signals for such processes?

Theoretical insight on these questions is presently rudimentary because the underlying processes are both dynamic (nonstatic) and nonperturbative. Progress
in the analysis of disoriented chiral condensate-generated physics and properties will be strongly driven by an F.A.D. experimental program.

We shall address issues (i) and (ii) below; for the moment, we turn to the question of experimental signals of disoriented chiral condensates.

2. A First Look at Experimental Signatures

We can use the physical picture of the late stages of the evolution of a disoriented chiral condensate developed in the previous section to develop a semi-quantitative understanding of some of the experimental signatures of such objects. We shall begin with an essentially classical picture, and then refine the picture to incorporate obvious quantum-mechanical effects.

We use the sigma model, without quarks, to model the chiral dynamics. In particular, we assume that the dynamical system can be described by the Lagrangian

\[ \mathcal{L} = \frac{1}{2} (\partial_\mu \bar{\pi} \partial^\mu \pi + \partial_\mu \sigma \partial^\mu \sigma) - \frac{\lambda}{2} (\bar{\pi}^2 + \sigma^2 - f_\pi^2)^2. \]

The minimum of the potential is at \( \sigma^2 + \bar{\pi}^2 = f_\pi^2 \), so that the chiral symmetry is spontaneously broken by having \( \langle \sigma \rangle = f_\pi \). Since \( \sigma \) is a scalar, this is the translation into the language of the effective theory of the belief that the scalar/isoscalar quark bilinear develops a vacuum expectation value in the full theory of QCD.

The dynamics we are interested in simplifies in the non-linear limit of \( \mathcal{L} \). That is, either \( \lambda \) is a Lagrange multiplier enforcing the constraint \( \bar{\pi}^2 + \sigma^2 = f_\pi^2 \), or, equivalently, \( \lambda \to \infty \). This non-linear limit seems to be a reasonable approximation: If one solves for the \( \sigma \) mass using the tree-level Lagrangian, one finds \( \lambda \approx \frac{m_\sigma^2 f_\pi^2}{(4 f_\pi^2)} \gg 1 \). One-loop corrections reduce this, but still leave \( \lambda \sim 6^{12} \). We have explicitly checked some aspects of our conclusions through numerical solutions at finite \( \lambda \); the conclusions are largely insensitive to this approximation.

With these assumptions we can model the relaxation process by studying the classical evolution of a domain-wall-like field configuration, suitably interpreting the eventual outgoing fields as quantum mechanical coherent states. The physical picture outlined above suggests that we assume spherical symmetry and begin our study of the dynamics at the time of hadronization of the spherical debris shell. Specifically, we assume that in the interior, the chiral condensate is disoriented by an angle \( \theta_0 \) from the \( \sigma \) direction, while in the exterior the chiral condensate is composed only of the \( \sigma \) field. Thus, it seems reasonable to assume that the fields
satisfy
\[ \bar{\pi}(r, 0) \to 0 \text{ as } r \to \infty, \]
\[ \bar{\pi}(r, 0) \to \hat{n}\pi \sin (\theta_0) \text{ as } r \to 0, \]
\[ \sigma(r, 0) \to f_{\pi} \text{ as } r \to \infty, \]
\[ \sigma(r, 0) \to f_{\pi} \cos (\theta_0) \text{ as } r \to 0. \]

In these equations \( \hat{n} \) is an arbitrary unit vector in isospin space.

Now, in order to study the dynamics, we need to make specific assumptions about the profile of both the fields and their velocities. We assume that the fields at time \( t = 0 \) can be modeled by a stationary kink-like configuration. It is easiest to express this by writing
\[ \bar{\pi} = f_{\pi} \sin \theta(r, t), \]
\[ \sigma = f_{\pi} \cos \theta(r, t), \]

with, for example,
\[ \theta(r, t = 0) = \theta_0(1 - \tanh \left( \frac{r - R_T}{T} \right)), \]
with \( \dot{\theta} = 0 \) and obvious candidate for a model of the interface.

There are several points to be made.

First, in this idealization of the dynamics, the evolution of the system is determined by an effective Lagrangian for \( \theta \), which is just
\[ \mathcal{L}_{\text{eff}} = \frac{f_{\pi}^2}{2} (\partial_{\mu} \theta)^2. \]

That is, the later stages of evolution are essentially that of a free field determined by the configuration of the disoriented condensate at the time of decoupling. This observation emphasizes that (in this scenario) the thickness of the interface must be determined by the thickness of the ‘hot’ shell of hadronic matter at the time of hadronization. This will permit us to make some semi-quantitative estimates of particle production.

Second, in this idealization (with a massless pion), the only contribution to the energy of the disoriented condensate comes from the interface between the interior and the exterior. We shall later generalize this to include the volume-dependent contribution arising from a non-vanishing pion mass.

How do we interpret this classical pion field? The obvious way of making a connection with quantum field theory is to interpret the outgoing wave as a
coherent state. Specifically, one might consider a coherent state

$$|\eta\rangle = e^{\int d^3k a^\dagger(k) \eta(k)} |0\rangle$$

with the property that

$$a(k)|\eta\rangle = \eta(k)|\eta\rangle,$$

where

$$\langle \eta|\eta\rangle = e^{\int d^3k |\eta(k)|^2}.$$  

The argument of the exponential is interpreted as the mean number of particles. (Note that we have temporarily suppressed isospin indices.)

There is, however, a subtlety associated with the interpretation in terms of coherent states which is extremely important both in the discussion of experimental signatures and in developing a deeper quantum-mechanical understanding of the disoriented chiral condensates. The issue involved is easily uncovered by the following series of considerations. First, we can consider a disoriented condensate in a specific cartesian isospin direction $\hat{n}$. If, in our classical discussion, both $\vec{\pi}$ and $\dot{\vec{\pi}}$ are in the $\hat{n}$ direction, then the field configuration has vanishing isospin density, and hence, vanishing total isospin. This is easily seen by noting that the isospin generators are given by

$$Q^a \sim \int d^3x \ \vec{\pi} \times \dot{\vec{\pi}}.$$  

However, the coherent states described above present us with a paradox. Consider a chiral condensate disoriented in, say, the $\pi_1$ direction. The coherent state will then be of the form

$$|\eta\rangle = e^{\int d^3k a^\dagger_1(k) \eta_1(k)} |0\rangle.$$  

But $a^{\dagger}_1 = 1/\sqrt{2} (a^{\dagger}_+ + a^{\dagger}_-)$, so that the coherent state will, for instance, have non-vanishing matrix elements with states of arbitrarily large charge!

The resolution to this is straightforward. All directions $\hat{n}$ are equally accessible, and so quantum mechanically one should average over them. This is equivalent to making a projection on the isospin-zero sector of the coherent states described above. This in turn leads us to the most striking experimental signal of the production of disoriented chiral condensates: a distribution of the fraction $f$ of neutral pions $\sim 1/\sqrt{f}$.  

We now give a simple derivation of the predicted isospin distribution of events with large number of pions, under the assumptions that the state is an isospin singlet, and that all of the pions have the same spatial wave function. The formulae
we shall derive are valid for any number of pions in such a state. In the limit that
the number of pions is large, we shall find that the distribution expected is \( \sim 1/\sqrt{f} \),
where \( f \) is the fraction of neutral pions.

Since we are concerned with \( 2N \) identical pions, we can carry out the derivation by introducing the operators \( a_i^\dagger, a_i \) which create/annihilate pions in the relevant state. (The index ‘i’ is a cartesian isospin index). We take these operators to be normalized so that \([a_i, a_j^\dagger] = \delta_{ij}\). The isospin generators are then realized by
\( I_i = -i\epsilon_{ijk}a_j^\dagger a_k \). While classically we shall find it useful to work in the cartesian
isospin basis, quantum mechanically it is more convenient to work in a basis in which \( I_3 \) is diagonal, and introduce appropriate raising and lowering operators. It is thus more convenient to work in terms of the creation operators \( a_3, a_\pm^\dagger = 1/\sqrt{2}(a_1^\dagger \pm ia_2^\dagger) \) and their hermitian conjugates. Similarly, it is convenient to work
with the z-component of the isospin generators
\[
I_3 = a_+^\dagger a_+ - a_-^\dagger a_-
\]
and the corresponding isospin raising/lowering operators
\[
I_\pm = \pm \sqrt{2}(a_\pm^\dagger a_3 + a_3^\dagger a_\mp). \]

We now turn to the construction of a multi-pion state ‘coherent’ state \(|\psi\rangle\)
which is an isosinglet and an eigenstate of the total pion number operator. Requiring
\( I_3|\psi\rangle = 0 \) implies that \( n_+ = n_- \), and hence that the total number of pions is even, as is the number of neutral pions. Thus we can expand
\[
|\psi\rangle = \sum_{n=0}^{2N} C_n^{(N)} (a_3^\dagger a_+^\dagger a_-^\dagger)^{N-n} |0\rangle,
\]
where \( 2N \) is the total number of pions, \( 2n \) is the number of \( \pi^0 \)’s, \(|0\rangle\) is annihilated by
the \( a_i \), and the \( C_n^{(N)} \) are expansion coefficients to be determined by the requirement
that \( I_\pm |\psi\rangle = 0 \).

We can determine the expansion coefficients up to an overall normalization and
phase fairly straightforwardly. First, observe that requiring \( I_+ |\psi\rangle = 0 \) gives us \( N \)
linear equations for the \( N+1 \) unknown \( C_n^{(N)} \). (The \( I_- \) equation yields no additional
information.) These equations can be solved in terms of a single coefficient, say
\( C_0^{(N)} \).
We can also by inspection find an explicit form for a $2N$ pion isosinglet state as follows. The operator
\[ S^\dagger = 2a_\uparrow^\dagger a_\downarrow^\dagger - (a_3^\dagger)^2 \]
commutes with the isospin generators, $[I_i, S^\dagger] = 0$. Consequently, the state
\[ |\psi\rangle = C_0^{(N)}(2a_\uparrow^\dagger a_\downarrow^\dagger - (a_3^\dagger)^2)^N |0\rangle \]
is a $2N$ pion isosinglet, and, by the considerations of the preceding paragraph, must be the most general such state in which all of the pions have the same spatial wave function.

From this it is possible to determine the probability for seeing $2n$ neutral pions out of $2N$ total pions in such a state:
\[ P(n, N) = \frac{(N!)^22^{2N}(2n)!}{(2N+1)!n!2^n} \]
Using Stirling’s approximation, it is straightforward to demonstrate that $P(n, N) \sim 1/\sqrt{n/N}$ in the limit that both $n$ and $N$ become large. These results have also been previously found using other arguments \(^1\text{13–16}\).

3. Tentative Phenomenology

Having established the striking signature of the $1/\sqrt{f}$ distribution of neutral pions, it is appropriate to try to develop the expected phenomenology in a little more detail. We are hampered a bit in that we cannot do any calculations from first principles. Nevertheless, we can make a number of estimates which should be useful guides in searching for disoriented chiral condensates. Similar estimates have been presented by Bjorken \(^4\).

We begin, as before, from the physical picture developed in the introductory section. That is, we consider a spherical region of radius $R$ within which the chiral condensate is disoriented by an angle $\theta_0$ from the sigma direction. The interior is separated from the exterior by an interface of thickness $T$ which is determined by the dynamics of the hot shell of hadronic matter resulting from the initial collision. We first estimate the energy density of the configuration at the time of hadronization of the shell. In the interior, we assume that there is essentially no kinetic energy, but there is an energy density arising from the non-vanishing pion
mass:
\[ \mathcal{E}_V \approx \frac{1}{2} m_\pi^2 \pi^2 \approx \frac{1}{2} m_\pi^2 f_\pi^2 \sin^2(\theta_0). \]

Similarly, there will be an energy density associated with the interface between the interior and the exterior. The dominant contribution in this region will be that due to the gradients of the fields:
\[ \mathcal{E}_S \approx \frac{1}{2} \left( \sum_a (\nabla \pi_a)^2 + (\nabla \sigma)^2 \right) \approx 2 \frac{f_\pi^2}{T^2} \sin^2(\theta_0/2). \]

The corresponding total energies are thus
\[ E_V \approx \frac{2}{3} \pi R^3 m_\pi^2 f_\pi^2 \sin^2(\theta_0) \]
and
\[ E_S \approx 8 \pi \frac{R^2}{T} f_\pi^2 \sin^2(\theta_0/2). \]

To proceed further, we need to estimate some of the relevant parameters. The shell cannot be thinner, at the time of hadronization, than the pion radius \( r_\pi \approx 2/3 \text{ fm} \). We take this to be our best estimate for \( T \). Thus, if we assume a dense layer of pions of this thickness, we have a total of
\[ N_{nor} \approx \frac{4 R^2}{r_\pi^2}. \]

Note that it seems reasonable to assume that these “normal” pions have typical average transverse momentum \( \langle p_t \rangle \approx .5 \text{ GeV} \). We note that this model predicts that “normal” multiplicity should scale with the surface area of the sphere. There is some evidence for this from studies of pion interferometry. These appear to be broadly consistent with the picture we are assuming, though details of the comparison are complicated by details of the data analysis. It is not clear to us whether the data favors a larger or smaller value for \( T \).

Our estimate of \( T \approx 2/3 \text{ fm} \) implies that the pions associated with the surface energy will have characteristic momenta of the order \( p_S \sim 2/T \approx .6 \text{ GeV} \). Thus, this component of the pion spectrum will be relativistic. Hence, we estimate the number of pions associated with the interface energy as
\[ N_S \approx \frac{E_S}{p_S} \approx 4 \pi R^2 f_\pi^2 \sin^2(\theta_0/2). \]

Similarly, assuming that \( R >> T \), so that the physical picture outlined above is meaningful, the pions associated with the interior volume will be non-relativistic.
Hence the total number of pions associated with the interior condensates is

\[ N_V \approx \frac{E_V}{m_\pi} \approx \frac{2}{3} \pi f_\pi^2 m_\pi R^3 \sin^2 (\theta_0). \]

While we cannot estimate thresholds in the absence of a tractable dynamical model of the early stages, we can define criteria for our picture to be self consistent. Specifically, the picture that we have suggested implicitly assumes that the system be macroscopic, with \( R >> T \). Taking \( R \sim 5T \sim 3 \text{ fm} \) seems a criterion for a ‘threshold’ above which our picture may be sensible. Taking \( f_\pi = 93 \text{ MeV} \) we thus estimate at this ‘threshold’: \( N_{nor} \approx 83 \) pions, \( N_S \approx 24.5 \sin^2 (\theta_0/2) \) pions, and \( N_V \approx 8.6 \sin^2 (\theta_0) \) pions. Such an object would have an energy in its center of mass frame of approximately

\[ M^* \approx E_{nor} + E_S + E_V \approx 41.5 \text{ GeV} + 16.0 \sin^2 (\theta_0/2) \text{ GeV} + 1.3 \sin^2 (\theta_0) \text{ GeV}. \]

If we assume that \( \langle \sin^2 (\theta_0/2) \rangle = \langle \sin^2 (\theta_0/2) \rangle = 1/2 \), which is probably reasonable as long as the energy of the condensate is small compared to the energy of the normal shell, we have \( M^* \approx 50 \text{ GeV} \) carried by \( \sim 100 \) pions.

As the size of the object increases, the importance of the chiral condensate increases because of the volume contribution. However, the coefficient of the volume term is rather small, so that it will not be the dominant contribution until \( R \sim 28 \text{ fm} \), corresponding to pion multiplicities of \( \sim 4000! \).

These estimates raise one issue which may be extremely important for the phenomenology. In our discussion we have divided the system at the time of hadronization into three parts: a surface layer of ‘normal’ pions, an interior dis-oriented chiral condensate, and an interface between the interior condensate and the exterior vacuum whose thickness is controlled by the ‘normal’ pions. Both the interface and the interior condensate contribute to the production of coherent pions, with the distinction serving mainly to suggest the momentum distribution of the condensate. One needs to worry, however, that the distinction between the ‘normal’ pions (which scale like \( R^2 \)) and the interface contribution to the coherent pions (which also scales like \( R^2 \)) is overdrawn. Our estimates suggest that they have comparable momentum scales, the ratio of the number of coherent pions from the interface energy to the number of normal pions is estimated to be

\[ \frac{N_S}{N_{nor}} \approx \pi f_\pi^2 f_\pi^2 \sin^2 (\theta_0/2) \approx 0.3 \sin^2 (\theta_0/2). \]

While for small values of \( \theta \) the distinction is probably reasonable, one must worry that when \( \theta_0 \approx \pi \) the dynamics may be rather more complicated than we have
suggested. That is, rather than the hot ‘normal’ component providing a barrier between the interior and exterior until the time of hadronization, if the energy in the interface is large enough, the interface might plausibly play a role toward ‘polarizing’ the hadronization of the putative normal component.

This issue may become a little more clearly drawn in a slightly different point of view (due to Bill Walker) regarding the mechanism for the formation of the disoriented chiral condensates\(^4\). The picture is easiest to describe in the rest frame of one of the colliding particles. In this frame, the oncoming particle is essentially a black disk as far as colored objects are concerned. (This picture is essentially due to the explosive growth of small-x gluons at large momenta.) As such, it acts as a ‘vacuum cleaner’, sweeping away all colored degrees of freedom associated with the target. What is left behind is ‘nothing’, but not vacuum. Since it will be a color singlet, it will not become attached to the disk, but will drift behind it, separated by a rapidity gap from the rest of the collision debris. This picture is essentially an alternative description of the same initial stage of the collision which we described in the opening section. As such, we would expect the ‘nothing’ to be characterized by a total energy and momentum, and to expand outward in the fashion described above, with a thin shell of ‘nothing’ (hot partonic matter) separating the rapidly cooling interior from the exterior. But in this picture, it is clear that as the ‘nothing’ cools below the relevant transition temperature, it is appropriate to describe it in terms of the sigma model. Thus, the distinction between the condensate and the ‘normal’ component is perhaps rather more blurred than we have been assuming, particularly for events in which \(\theta_0 \approx \pi\).

Finally, we turn to the question of estimating cross sections for events such as we have been describing. This is exceedingly difficult in the absence of a more detailed understanding of the mechanism for the formation of the objects. Bjorken has suggested an approach to this question\(^4\). Roughly speaking, we know that in order to have a ‘thermalized’ state of rest energy \(M^*\), we need to have at least something of the order of \(M^*/m_0\) strongly interacting partons suitably localized. (\(m_0\) is a QCD scale parameter \(\sim 1\ GeV\).) Then the question becomes whether or not this is realized in hadron-hadron collisions of a given energy. This depends critically on the small-x behaviour of the structure functions; the QCD-inspired model of Block, Halzen and Margolis\(^{17}\) simulates this behavior, and indicates that the bound is probably easily satisfied at SSC energies. However, it is not yet clear that the energy is deposited in a suitable configuration. A simple geometric argument based on the constituent quark model suggests a lower bound of approximately a nanobarn, but this is probably rather conservative, at least at supercollider energies.
4. Are Centauros Related to Disoriented Chiral Condensates?

Bjorken has suggested that the Centauro phenomena might be interpreted in terms of disoriented chiral condensates. We now turn to this question.

Centauro events are cosmic ray events exhibiting:

- Large (∼ 100) numbers of hadrons;
- Little apparent electromagnetic energy and hence, no π⁰'s;
- High hadronic $p_t$, reported as $k_\gamma(p_t) = 0.35 \pm 0.15 \text{ GeV}$, where $k_\gamma$ is the photon elasticity;

In addition to the Centauro events, a class of hadron-enriched events has also been reported. Cosmic ray events with $\sum E_{\text{tot}} \geq 100 \text{ TeV}$ are presented on a scatter plot of the number of hadrons $N_h$ versus the fraction $Q_h = \sum E_\gamma^h / (\sum E_\gamma^h + \sum E_\gamma)$ of the visible energy which these hadrons constitute. When compared with Monte-Carlo simulations of families based on models of the strong interactions, and assuming that cosmic ray primaries are predominantly protons, there are far too many (∼ 20%) events in regions not populated by the Monte-Carlo. These events show fluctuations in hadron number and/or energy fraction.

Before continuing it is probably necessary to briefly review the status of candidate Centauro events. The 5 ‘classic’ Centauro events were seen in the two-storeyed emulsion chamber experiment of the Brasil-Japan collaboration, located at 5220 m at the Chacaltaya observatory in Bolivia. At least two additional candidate Centauro events have also been seen in Chacaltaya chambers. At least one additional candidate has been seen in the Pamir emulsion chamber experiment. On the other hand, it is claimed that Centauros have not been observed in emulsion chambers at Mt. Kanbala (5500 m, China-Japan Collaboration) or at Mt. Fuji (3750 m., Mt. Fuji Collaboration), despite comparable cumulative exposures. More precisely, the China-Japan collaboration reports an upper limit of the fraction of such events among hadron families with energy greater than 100 TeV to be 3% at the 95% confidence level. This appears to be a limit incompatible with the rate at which Chacatalya has observed Centauros. This comparison may be too glib, however, because of differences in emulsion chamber design and data analysis. Of particular importance may be the differing techniques for separating hadronic showers from others.

The two groups are similarly divided on the (non)observation of the more general class of hadron-enriched events mentioned above. In this context, the debate (which has been going on for over a decade) seems to reduce to a question of whether the data signal a change in the composition of cosmic ray primaries at these energies, or whether they signal a change in the hadronic interactions.
With these caveats, we proceed to discuss the consistency of the Centauro and hadron-enriched phenomena with the phenomenology of chiral condensates as developed in the preceding sections.

We begin with the characteristic feature of Centauro’s and the hadron-enriched events: anomalously large amounts of energy in the hadron component. The suppression of $\pi^0$’s which this implies is often taken to indicate a suppression of pions altogether. The argument is basically statistical: one would ordinarily expect the neutral fraction to be given by essentially a binomial distribution, resulting in events sharply peaked about $1/3$. As we have seen above, however, the distribution for an isospin-zero coherent state of pions is $\sim 1/\sqrt{f}$, with $f \sim 0$ being the most probable fraction. (Note, however, that $\langle f \rangle = 1/3$.) Thus, one is tempted to interpret the classic Centauro events as signals of a disoriented chiral condensate.

A problem immediately arises. While the multiplicity of the Centauro events is comparable to our estimates, we also suggested that a large fraction of the hadrons would be ‘normal.’ This would appear to rule out Centauro events as signals of a disoriented condensate. The only apparent resolution to this difficulty in interpretation is that these events represent a situation for which our understanding of the interface between the interior and the exterior at time of hadronization is suspect: $\theta_0 \approx \pi$.

Assuming this to be true, we are faced with an additional problem: the claimed anomalously high $\langle p_t \rangle$ of the Centauros. This value depends on the fact that Centauro I was close enough to the detector that the position of the interaction vertex could be determined by the angular divergence of the showers in the detector. The actual value measured is

$$\frac{\langle E_{h}^{(\gamma)} R_h \rangle}{H} = k_{\gamma} \langle p_t \rangle = 0.35 \pm 0.15 \text{ GeV},$$

where $E_{h}^{\gamma}$ is the portion of an incident hadron energy which is converted to (visible) electromagnetic energy, $R_h$ is the distance of an incident hadron from the center of the event, and $H$ is the height of production. In order to determine $p_t$, however, one needs to know the value of the gamma-ray inelasticity, $k_{\gamma}$. The range is usually quoted as $0.2 - 0.4$, with the lower range being preferred for nucleons, while the higher end is preferred for pions. Direct measurement of $k_{\gamma}$ in emulsion chambers is impossible because of the high energy threshold ($\sim 1 \text{ TeV}$). As a result, estimates are based on extrapolating accelerator data or on Monte-Carlo simulations. These seem to indicate that one should use $k_{\gamma} \sim 0.4$ or larger\textsuperscript{21}. We would thus estimate that $\langle p_t \rangle \sim 0.875 \pm 0.375 \text{ GeV}$ with large systematic uncertainty. This seems compatible with our estimates of the previous section. We note most analyses
have followed the Japan-Brazil collaboration and have used $k_\gamma = 0.2$, based on the assumption that the hadrons are nucleons.

We seem to be on somewhat firmer ground in confronting the hadron-enriched events. Detailed comparison will require rather careful Monte-Carlo simulations, but the general features of the plot of $N_h$ versus $Q_h$ seems to be entirely consistent with a phenomenology based on disoriented condensates, with corresponding fluctuations in the neutral fraction, and in $\theta_0$.

We are thus rather encouraged that the Centauro and the hadron-enriched events may be signals of the creation of disoriented chiral condensates.

Finally, if we accept this interpretation, then we can use the experimental data to suggest both the cross section and the threshold energy for the production of disordered chiral condensates. As far as the cross section is concerned, the abundance of the hadron-enriched events suggest that they are almost generic at these energies (cosmic ray showers with $> 100$ TeV of visible energy). We were unable to track down more detailed estimates of the energies in the typical hadron-enriched events, but one can determine a lower bound on the threshold energy for Centauros. In the 5 classic events, the estimated visible energy of the showers (after correcting for atmospheric effects) was $\sim 350$ TeV. Dividing by $k_\gamma \sim 0.4$ yields an estimated total energy of 875 TeV. If the parent interaction is due to a collision between nucleons, then we can estimate the minimum center of mass energy of the collision as $\sqrt{s} \sim 1.3$ TeV. This suggests that there is no inconsistency with the negative results of the UA1 and UA5 Centauro searches\textsuperscript{22–24}.

5. Experimental Challenges

We conclude that there is a considerable body of evidence based on high-energy cosmic-ray observations, buttressed by some theoretical considerations, that there may be a significant fraction of the pp and p\bar{p} inelastic cross section above $\sqrt{s} \sim 1.3$ TeV that have ‘anomalous’ pionic multiplicity patterns. We have described one model for accounting for such multiplicity patterns above a certain threshold energy. The Tevatron energy is above this threshold and lies in the lower range of the spectrum of anomalous cosmic-ray events; the LHC and the SSC are well above the anticipated threshold energy and cover the spectrum of most of the cosmic-ray events.

The characteristics of these events and their interpretation have been a subject of widespread interest in the cosmic-ray, particle-physics, and astrophysics communities for over a decade. This continued interest represents considerable justification for continuing their experimental investigation at the Tevatron, the LHC, and the
SSC. The replication of the below-threshold UA1 and UA5 minimum-bias experiments at the Tevatron would constitute a natural bridge for similar experiments at super-collider energies and would replace the lower-energy cosmic-ray data, whose sample size is very limited and whose processing is both complicated and model dependent, by much more extensive experimental data not subject to the same procedural uncertainties. These data would provide extremely valuable guidance for the choice and refinement of the theoretical modeling of the production and symmetry-breaking mechanisms of strong-interaction physics that involve complicated collective effects rather than better understood perturbative processes. This guidance is likely to be more unequivocal, at least in regard to chiral-symmetry breaking mechanisms, than corresponding investigations at RHIC where large nuclear/deconfinement effects are likely to lead to much more complicated signals.

In summary, collider experiments at the Tevatron, the LHC, and the SSC looking for generic, high-threshold, unusual-multiplicity-pattern hadron-hadron inelastic collisions are needed for the following reasons:

1. The intrinsic interest of the continued investigation of unusual strong interaction production mechanisms at Tevatron (and higher) energies.

2. To replace most of the limited-sample-size cosmic-ray data set (which is based on a variety of observational setups, models, and assumptions) by the results of controlled experiments.

3. To provide the experimental input that is needed to discriminate between the large number of theoretical conjectures that have been proposed to account for the cosmic-ray results as well as those that relate to the symmetry-breaking mechanisms of QCD. Experimental guidance is urgently needed for the modeling of the complicated collective effects that seem to be responsible for physics of this kind.
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