COHERENCE AND DECOHERENCE IN RADIATION OFF COLLIDING HEAVY IONS

André KRZYWICKI

Laboratoire de Physique Théorique et Hautes Energies, Bât. 211, Université de Paris-Sud, 91405 Orsay, France

Abstract

We discuss the kinetics of a disoriented chiral condensate, treated as an open quantum system. We suggest that the problem is analogous to that of a damped harmonic oscillator. Master equations are used to establish a hierarchy of relevant time scales. Some phenomenological consequences are briefly outlined.

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1 INTRODUCTION

It is plausible that, in high-energy collisions, one occasionally produces a bubble of classical pion field, randomly oriented in isospace [1, 2, 3]. There is a growing literature on this subject [4]-[10]. All these admittedly very speculative papers are not purely esoteric, however. A transient production of a ”disoriented chiral condensate” could be an explanation for the heretofore mysterious Centauro events reported by cosmic ray people [11]. Indeed, it is easy to convince oneself that the fraction \( r \) of neutral pions produced in the decay of a bubble of disoriented chiral condensate should be distributed according to the probability law

\[
dP(r) = \frac{1}{2\sqrt{r}} dr
\]

Hence, the probability of events with small \( r \) is unusually high compared to naive statistical expectations.

The specific scenarios proposed in the theoretical papers quoted above are more complementary than equivalent. Hopefully new data will select the right one. For the time being, it seems reasonable to develop all the a priori plausible scenarios, checking their self-consistency and looking for implications. Thus, the aim of this paper is to further develop the ideas put forward in ref. [3]. We follow a strategy similar to that adopted there, sacrificing generality in favour of simplicity.

The plan of the paper is the following: In sect. 2 we recall some of the results obtained in ref. [3] in 1+1 dimensions and we discuss a possible extension to 1+3 dimensions. We argue that the disoriented chiral condensate should be considered as an open quantum system. In sect. 3 we formulate the analogy between our problem and the kinetics of a laser under threshold. In sect. 4 we establish a hierarchy of time scales relevant to the evolution of the condensate. Sect. 5 contains the discussion.

2 THE ORIGINAL MODEL AND GENERALISATIONS

The starting idea is to use the classical equations of motion derived from an effective chiral Lagrangian to calculate the intensity of soft-pion radiation in a
high-energy heavy-ion collision. In ref. \[3\] the lowest order chiral Lagrangian, viz. that of the non-linear $\sigma$ model has been employed:

$$L = \frac{1}{2} f^2 \left[ (\partial \sigma)^2 + (\partial \pi)^2 \right] \quad (2)$$

and

$$\sigma^2 + \pi^2 = 1 \quad (3)$$

The classical equations of motion are summarized in the following current conservation equations

$$\partial V = 0 \quad \partial A = 0 \quad (4)$$

where $V_\mu$ and $A_\mu$ are the Noether isovector and isoaxial currents associated with the global $SU(2) \times SU(2)$ symmetry of the model. These equations must be supplemented by appropriate boundary conditions. Inspired by Heisenberg’s old papers \[12\] on semi-classical description of multi-particle production, we have adopted the following idealization: At time $t = 0$ the whole energy of the collision is localized within an infinitesimally thin slab with infinite transverse extent (instead of a pancake shaped region). The symmetry of the problem then implies that the pion field depends only on the invariant $s = t^2 - x^2$, where $x$ is the longitudinal coordinate. Thus the problem becomes effectively 1 + 1 dimensional and the currents take the form

$$V_\mu = a \, \partial_\mu \phi \quad A_\mu = b \, \partial_\mu \phi \quad (5)$$

where $a, b$ are integration constants satisfying $a \cdot b = 0$ and

$$\partial^2 \phi = 0, \ s > 0 \quad (6)$$

Hence, the problem is reduced to an Abelian one. The solution of (6) corresponding to currents non-vanishing within the light-cone is $\phi(s) = \log(s/s_0)$. 

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Projecting the pion field on the axes of the triad $a, b, c = a \times b$ and after some algebra one finds for $s > 0$

$$\begin{align*}
\pi_a &= 0 \\
\pi_b &= -\sin(\kappa \phi) \\
\pi_c &= \left(\frac{a}{\kappa}\right) \cos(\kappa \phi)
\end{align*}$$

with $\kappa = \sqrt{a^2 + b^2}$. Of course, the classical solution breaks the symmetry of the theory: within the light-cone the pion field rotates in isospace in a plane perpendicular to a randomly chosen vector $a$. Outside the light-cone $\sigma = 1$. The solution is singular for $s = 0$. Hence, it describes radiation by sources living on the light-cone. For $t = const$, the pion field fills the space between these two sources and oscillates violently as one approaches them. One then enters the region where is stored almost all the energy of the collision and where the gradient expansion involved in the derivation of the effective Lagrangian from QCD stops making sense. Close to the sources the $\sigma$ model (1) cannot be trusted and, in our opinion, its use is not justified.

Thus, the picture in 1+1 dimensions is essentially the same as in [2]: there is an inner region, where resides a disoriented chiral condensate, separated by a potential barrier from the outer physical vacuum. All this comes out quite naturally. Unfortunately, the generalisation to 1+3 dimensions is not unique. One can assume as in refs. [2, 6] that somehow the same topology is created also in this case. This occurs perhaps in events with exceptionally high multiplicity. However, in this paper we shall keep working with the boost-invariant scenario [12, 13] as in [3], postulating that the sources are localized in those space regions where energy density is high in a generic heavy-ion collision (i.e. essentially within the two receding pancakes). We shall see soon that this assumption is less innocent than it might appear and leads to difficulties in the purely classical framework. The way out will be to consider the disoriented chiral condensate as an open quantum system, which it actually is.

Before going further we need to consider for a while the modifications of the results obtained in 1+1 dimensions, when the initial system has a finite space extent $\epsilon$. As in [3] we shall calculate the total radiated energy at some large time $t$, but we shall not include in the space integration the regions where reside the sources. Technically, in calculating the spectrum,
we Fourier transform the fields multiplied by a cut-off function \( f(t, x) \) which equals 1 within the light-cone, except in the regions \(|x| \leq |t| - \epsilon\) where it vanishes rapidly. In analogy to [3] we are led to calculate the integral of the type

\[
I = \int_0^t dx f(t - x)e^{ikx + \alpha \ln (t^2 - x^2)}, \quad k > 0
\]  

(8)

(for simplicity of writing we consider the integration over \( x > 0 \)). Change the integration variable to \( z = t - x \) and consider the integration over the closed contour made up of three parts: \( 0 < \text{Re} z < t, -\frac{\pi}{2} < \arg z < 0 \) and \(-t < \text{Im} z < 0 \). For large \( t \) one has

\[
I \simeq -i(-2it)^\alpha e^{ikt} \int_0^t dy f(-iy)e^{-ky}y^\alpha[1 + O(y/t)]
\]  

(9)

When \( f = 1 \) one can obtain from (8) the leading term of the large-\( t \) expansion of the exact expression quoted in [3] (for \( t \to \infty \) the integral converges towards \( \Gamma(1 + \alpha)/k^{1+\alpha} \)). When \( \epsilon > 0 \) (and \( f(z) \) is smooth enough, e.g. \( f(z) \sim \exp(-\epsilon/|z|) \)) the contribution to the integral in (8) from the region \( 0 < y \lesssim \epsilon \) is vanishing and there appears a factor \( \exp(-\epsilon |k|) \) in the spectrum. This agrees with the physical intuition: only modes with wavelength larger than the size of the emitter are effectively produced. One could arrive at the same conclusion checking that the stationary point of the exponential in the integrand of (8) is at a distance of order \( O(1/k) \) from the light-cone and is therefore outside of the cut-off region provided \( k \) is small enough.

In 1 + 3 dimensions we have been unable to find the most general solution of eqs. (4) for some physically reasonable set of boundary conditions. However, if one accepts to limit oneself to solutions of the form (7), then the problem again becomes Abelian and the field equations reduce to (6) (outside of sources). We have studied various solutions of the wave equation (6). The upshot of this study is almost evident. Therefore, we shall not enter into details, limiting ourselves to pointing out what is truly relevant. The solutions of interest are those with cylindrical symmetry : \( \phi = \phi(s, r) \), where \( r \) is the transverse coordinate. The initial transverse radius of the source is denoted by \( R \). It is not obvious whether \( R \) should be assumed to be roughly equal to the nuclear radius of the colliding heavy ions or should rather be set to a smaller value. It depends on how much coherence one expects to have in the transverse direction and this in turn depends on the mechanism
triggering the creation of the disoriented chiral condensate. Most likely, $R$ fluctuates from event to event and there is a penalty for $R$ being large (in nuclear units). The generic behaviour of $\phi$ is

$$\phi \sim \log s, \sqrt{s}, r \lesssim R$$

and

$$\phi \sim s^{-1}, \sqrt{s} \gg R, r$$

Eq. (10) simply states that the one-dimensional idealization of ref. [3] is roughly realistic as long as the sources are separated by a distance less than $R$. Eq. (11) asserts that very far from a finite size source the field becomes that of a point emitter. Notice, that as times goes on the bulk of the field stays at a finite invariant distance from the sources. In the analogue of eq. (3) the stationary point is at a distance from the source, which inevitably becomes less than $\epsilon$ provided $t$ is large enough and whatever is $k$. In other words for $t \to \infty$ the field collapses on the sources and there is no radiation. A straightforward generalisation of the discussion of ref. [3] simply does not work. Of course, it is completely unrealistic to imagine that the two sources move indefinitely intact and with the velocity close to that of light. However, we do not see any reason why they would switch off suddenly at $t \simeq R$, before the radiation enters into the three-dimensional regime.

Trying to understand the physics of this problem, one should notice that the soft-pion radiation we are considering differs in (at least) one essential aspect from the familiar bremsstrahlung from a classical current discussed in any elementary textbook of quantum electrodynamics. The point is that the pion system is permanently in contact with the debris of the colliding nuclei. Thus the pion system is an open one and is subject to decoherence. We shall see in the next sections under which conditions the latter can occur before the one-dimensional expansion is over.

3 THE DAMPED OSCILLATOR ANALOGY

We shall focus in this section on the decoherence problem. Because the subject is perhaps not very familiar to all the potential readers of this paper, and also in order to spell out the numerous simplifications we are going to
adopt, we shall briefly sketch certain derivations. For more details and for
references to the original works the reader should consult, for example, refs. [14, 15, 16].

The idea is to consider the disoriented chiral condensate as a ”system” in
contact with thermalized hadronic matter produced in a heavy-ions collision
(the ”bath”). We shall represent the ”bath” by a collective variable $B$, which
we shall couple linearly to the ”system”. One can write $B = \langle B \rangle + \delta B$.
The classical component of $B$ will yield sources responsible for the soft-pion
radiation discussed in [3]. The quantum fluctuations of $\delta B$ will tend to
break the coherence of the radiation. Actually, in order to avoid technical
complications we shall consider an Abelian toy model in 1+1 dimensions. The
non-linear $\sigma$ model is reducible to this model under circumstances mentioned
in the preceding section, but in the classical limit only. However, we hope that
the toy model is sufficient to produce rough estimates of the characteristic
time scales we are interested in.

The starting point is the von Neumann-Liouville equation satisfied by the
density matrix, denoted by $W$:

$$\dot{W}(t) = -i[H, W(t)]$$  (12)

where $H$ is the Hamiltonian describing a free scalar field coupled linearly to
the collective variable representing the ”bath”. We assume that the coupling
(in the Lagrangian) is a gradient one, viz. $\sim \partial_\mu \phi \partial^\mu B$, so that only a rapidly
varying ”system” field is strongly coupled.

We partition the rapidity space into cells of extent $\delta y > 1$. We denote
by $a^\dagger$ the operator creating a quantum of our scalar field within a given cell
and we write

$$H = \omega a^\dagger a + H_{bath} + V$$  (13)

with the interaction

$$V = (a + a^\dagger)j(t) + \lambda(\omega)[ab^\dagger + a^\dagger b]$$  (14)

Here, $\omega$ is the average energy within a cell while $j(t)$ and $b$ are the appropriate
projections of the classical current $j \sim \langle \partial^2 B \rangle$ and of the quantum ”bath”
field $\delta B$, respectively. We assume that $H_{bath}$ does not couple distinct cells.
Strictly speaking, the time resolution $\delta t$ is finite and the time derivatives in the
following text are to be understood as the coarse-grained rates of change of
the corresponding quantities. However, we shall use the continuum notation
for simplicity.

Going over to the interaction representation we get

\[ \dot{W}_I(t) = -i[V_I(t), W_I(t)] \quad (15) \]

We assume that \( W(t) \) factorizes at \( t = 0 \)

\[ W(0) = \rho(0)X \quad (16) \]

where \( X \) is a stationary density matrix describing the internal dynamics of
the "bath". We again make an idealization: of course, the "bath" is cooling
as time goes on. However, this cooling is expected to be slow during the
one-dimensional regime we are interested in and we neglect it altogether (in
ref. [13] it is found, using the hydrodynamical model, that the temperature
falls with proper time like \( \tau^{-\frac{1}{3}} \) only). Furthermore, the "bath" is not at rest
with respect to the "system". To deal with this complication we adopt a
simple ansatz:

\[ X \sim \exp[-\frac{1}{2} \beta P_{\mu}(u_1^\mu + u_2^\mu)] \quad (17) \]

where \( P_{\mu} \) is the energy-momentum operator of the "bath" and \( u_i \) are covariant
velocities of the outward expansion of the "bath". For \( u_1 = u_2 = (1, 0) \) \((17)\)
becomes the standard canonical density. We assume for definiteness that
\( u_1^1 = -u_2^1 \) in the rest frame of the condensate. Then the exponent becomes
\[ -\beta \gamma H_{\text{bath}} \].

The spectrum of the "bath" is red-shifted due to its expansion
and has an effective temperature \( T/\gamma \). The Lorentz factor \( \gamma \) is merely a
phenomenological parameter.

For \( \lambda = 0 \) the problem has the well known solution \([17]\)

\[ W_I^{(0)}(t) = D[\alpha(t)]|0\rangle\langle 0|D^\dagger[\alpha(t)]\}X \quad (18) \]

where

\[ \alpha(t) = i \int_0^t dt' j(t')e^{i\omega t'} \quad (19) \]

is the Fourier component of the classical solution of the equations of motion
and \( D(\alpha) \) is the unitary operator.
\[ D(\alpha) = \exp[\hat{\alpha}a - \alpha a^\dagger] \] (20)

Defining
\[ \tilde{W}(t) = D^\dagger[\alpha(t)]W_I(t)D[\alpha(t)] \] (21)

and similarly for other operators, we reduce the von Neumann-Liouville equation to
\[ \dot{\tilde{W}}(t) = -i[\tilde{V}(t), \tilde{W}(t)] \] (22)

The perturbative solution of (22) is
\[ \tilde{W}(t) = \tilde{W}(0) + \sum_{n=1}^{\infty} (-i)^n \int_0^t dt_1 \cdots \int_0^{t_{n-1}} dt_n [\tilde{V}(t_1), \ldots, [\tilde{V}(t_n), \tilde{W}(0)] \ldots] \] (23)

We work to order \( O(\lambda^2) \) only and take a trace over the "bath" degrees of freedom in order to derive a master equation for the reduced density matrix \( \tilde{\rho}(t) = \text{tr}_{\text{bath}} \tilde{W}(t) \). The perturbative approach is, of course, meaningful only when the interaction between the "bath" and the "system" is weak enough not to upset the assumed thermal equilibrium of the "bath". The derivation starts from the observation that \( \tilde{\rho}(t) \) is obtained from \( \tilde{\rho}(0) \) by the action of a linear operator, call it \( U(t) \). A simple result is obtained in the so-called Markovian approximation, viz. replacing the operator \( \dot{U}(t)U^{-1}(t) \) by its value at \( t = \infty \). The main conditions for this to be realistic is that the relaxation time of the "bath" is much shorter than that of the "system". We adhere to this idealization, although in our problem it is likely to be particularly crude, since "bath" oscillations appear slowed down due to its expansion.

Using the notation
\[ T_{\text{bath}}(\ldots X) \equiv \langle \ldots \rangle \] (24)

we have
\[ \langle b(t) \rangle = 0 \] (25)

since \( b \) is linear in \( \delta B \). We also assume, for the sake of simplicity, that
\[ \langle b(t)b(t') \rangle = 0 \] (26)

The master equation is then

\[ \dot{\rho} = -i\lambda^2 \Delta [\bar{a}^\dagger a, \bar{\rho}] + \lambda^2 \kappa \{ [\bar{a}, \bar{\rho}^\dagger] + [\bar{a}\bar{\rho}, \bar{a}^\dagger] \} + 2\kappa n \lambda^2 [\bar{a}, [\bar{\rho}, \bar{a}^\dagger]] \] (27)

where

\[ \kappa + i\Delta = \int_0^\infty dt e^{i\omega t} \langle [b(t), b^\dagger(0)] \rangle \] (28)

\[ \kappa n = \text{Re} \int_0^\infty dt e^{i\omega t} \langle b^\dagger(0)b(t) \rangle \] (29)

Eq. (27) is formally identical to that governing the behaviour of a damped quantum oscillator. Eq. (19) implies that generically \( \alpha(t) \) becomes independent of time provided \( t > t_0 \), with the rough estimate

\[ t_0 \omega \sim 1 \] (30)

Thus, for \( t > t_0 \) we can drop the tildas in (27). In this way we finally reduce our problem to that of a damped oscillator. Of course, we have by no means demonstrated that the kinetics of a disoriented chiral condensate is equivalent to that of a damped quantum oscillator. We have postulated this analogy from the outset, because we feel it may give some insight into the real problem. The aim of this section is only to explain the real significance of the analogy, by making explicit the dynamical postulates it involves.

4 THE HIERARCHY OF TIME SCALES

Let us go to the Schroedinger representation and let us write the reduced density matrix in the Glauber representation

\[ \rho(t) = \int d^2 \beta |\beta\rangle P(\beta, \bar{\beta}, t) \langle \beta| \] (31)

It is well known that a master equation for \( \rho \) can be converted into a Fokker-Planck equation for the Glauber function \( P \). In the case of a damped oscillator the solution of the latter equation is known explicitly:
\[
P(\beta, \bar{\beta}, t) = \int d^2 \xi \, P(\beta, \bar{\beta}, t|\xi, \bar{\xi})P(\xi, \bar{\xi}, t_0)
\]  
(32)

where
\[
P(\beta, \bar{\beta}, t|\xi, \bar{\xi}) = \frac{\eta}{\pi} e^{\eta|\beta - \xi e^{-\theta(t-t_0)}|^2}
\]  
(33)

and
\[
\eta^{-1} = n[1 - e^{-2\lambda^2\kappa(t-t_0)}]
\]
\[
\theta = i(\omega + \lambda^2 \Delta) + \lambda^2 \kappa
\]  
(34)

Assume, that \(\lambda\) is so small that at \(t = t_0\) the ”system” is still almost coherent, viz. \(P(\xi, \bar{\xi}, t_0) \simeq \delta^2(\xi - \alpha(t_0))\). Then the Glauber function at \(t > t_0\) is given by the right-hand side of (33) with \(\xi = \alpha(t_0)\). It is an easy exercise to calculate the elements of the density matrix in the occupation number representation. One finds
\[
\langle N|\rho|N \rangle = (n + 1)^{-1} \left[ \frac{n}{n + 1} \right]^N [1 + O(e^{-2\lambda^2\kappa t})]
\]  
(35)

and
\[
\langle N|\rho|M \rangle \sim e^{-\lambda^2\kappa|N-M|t}, \, N \neq M
\]  
(36)

We are now in position to identify the different time scales entering our problem. First, there is the time \(t_0 \gtrsim \omega^{-1}\) required to build up the bremsstrahlung field. For a given rapidity cell, the probability of finding \(N\) quanta is Poissonian with average \(\bar{N} \simeq \alpha(t_0)^2\) and width \(\sqrt{\bar{N}}\). As time goes on the off-diagonal terms of the density matrix die out following the law (36). The coherence between the relevant distinct multiplicity states of the bremsstrahlung field is lost for \(t > t_{decoh}\), with
\[
t_{decoh} \simeq [\lambda^2\kappa\sqrt{\bar{N}}]^{-1}
\]  
(37)

Finally, the ”system” reaches a stationary state. Indeed, the fluctuation-dissipation theorem implies that \(n^{-1} = \exp(\beta\omega) - 1\). Hence, it follows from eqs. (35)-(36) that for \(2\lambda^2\kappa t >> 1\) the density matrix is \(\langle N|\rho|M \rangle \sim \delta_{MN} \exp(-\beta\omega N)\). The thermalization time-scale \(t_{therm}\) is clearly
\[ t_{\text{therm}} \simeq [\lambda^2 \kappa]^{-1} \]  

(38)

For large enough multiplicity and small coupling between the "system" and the "bath" one finds the following hierarchy of time scales

\[ \omega^{-1} < t_0 < t_{\text{decoh}} << t_{\text{therm}} \]  

(39)

The second inequality is a dynamical constraint. It should be satisfied if the transient creation of a coherent classical field is to take place.

The estimates (37) - (39) have so simple an intuitive meaning that one is tempted to believe that their validity transcends the over-simplified model used to get them.

5 SUMMARY AND DISCUSSION

The foregoing discussion has actually two facets. We have started by trying to extend the results of [3] to 1+3 dimensions. In this attempt we have encountered difficulties, which are at least partly rooted in the assumptions we have adopted. Of course, these assumptions may just reflect our prejudices. Anyhow, we have noticed that the difficulties can be circumvented if one considers the disoriented chiral condensate as an open quantum system. Hence, we have embarked into the discussion of the time scales entering the kinetics of the condensate, insisting on the (presumed) analogy with the damped quantum oscillator (or laser under threshold). However, the relevance of discussing decoherence does not depend on the soundness of our original motivation. This question has to be faced in all scenarios. Therefore, consider it first.

As one might expect, the decoherence time scale depends on the (Fourier transform of) the linear response function of the "bath" and on the strength of the coupling between the "system" and the "bath". A thorough discussion of these issues goes beyond the scope of this paper. In the present state of the art it would inevitably involve much model building. Let us limit ourselves to some rather obvious comments and guesses.

From the definition of the spectral function \( \kappa(\omega) \) it follows that it satisfies the sum rule
\int d\omega \kappa(\omega) = 1 \quad (40)

Let \( \Gamma \) denote the size of the bandwidth of the "bath". Roughly speaking \( \kappa(\omega) \simeq \Gamma^{-1} \) for \( \omega \) within the bandwidth and zero otherwise. Remember, that \( \kappa(\omega) \) has to be calculated setting the temperature of the "bath" to the effective value \( T/\gamma \). In a very high-energy collision this effective temperature is presumably low and it is perhaps reasonable to guess that the bandwidth extends from the lowest frequencies to \( \Gamma \), and that the latter is controlled by a hadronic scale, say \( \Gamma \lesssim 1 \text{ GeV} \).

In the relativistic regime, the assumed gradient nature of the coupling between our "system" and the "bath" implies that the coupling \( \lambda = \lambda_0 \omega \), where \( \lambda_0 \) is a dimensionless constant. Since the terms in the action involving the field \( B \) are supposed to mimic the interactions neglected when one keeps only the lowest order term in the chiral action, it is perhaps reasonable to guess that \( \lambda_0 \simeq f_\pi/M \), where \( M \) is the mass scale where the chiral perturbation theory breaks down (\( M \simeq 1 \text{ GeV} \), say). Considering a heavy-ion collision one should further multiply the coupling term, i.e. the classical current \( j \) and the coupling \( \lambda \), by the dimensionless constant \( f_\pi R \). One then finds \( \bar{N} \propto R^2 \), as it should be.

An extension of the discussion to the non-Abelian case presents additional difficulties. However, as long as chiral-symmetry is maintained, the disorientedness of the chiral condensate should not be affected by the decoherence.

Notice, that at \( t \simeq t_{\text{decoh}} \) the multiplicity distribution is still nearly Poissonian. The coherence between different multiplicity states has been broken, but correlations had no time to build up. However, for \( t > t_{\text{decoh}} \) the evolution of the system is no longer governed by the classical equations of motion.

Let us now return to the scenario outlined in sect. 2. As argued there, the classical field should loose its coherence before the end of the one-dimensional regime. This would invalidate the argument based on the classical equations of motion and maintaining that the pion field will collapse on the sources during the three-dimensional expansion. The condition for that not to happen is

\[ t_{\text{decoh}} \lesssim R \quad (41) \]

This is a stringent condition. It implies that only the quanta belonging to the bandwidth of the "bath" may contribute to the observable signature of
the disoriented condensate. Thus, the decay products of the condensate are expected to be localized in rapidity, presumably within an interval a few units long. This is to be contrasted with the purely classical result of ref. \cite{3}, where a rapidity plateau was found.

If our guess for $\Gamma$ and $\lambda_0$ is not too far from reality, the condensate should decay into a large number of pions, of the order of $10^2$ for $R \approx 3$ fm. Otherwise, the constraint (41) would be hard to meet. An extension of the model to $R \lesssim 1$ fm is problematic. If it is nevertheless attempted, then for $\lambda_0 \simeq 0.1$ the satisfaction of (41) requires either an unreasonably high multiplicity or a very narrow bandwidth. This probably means that one should not expect to discover the disoriented chiral condensate in hadron-hadron collisions when final states have generic overall topology. Here we join the intuition of the authors of refs. \cite{2,6}.

We would like to end this very speculative paper with a few general remarks:

- Contrary to what we have stated at the end of ref. \cite{3}, describing the creation and evolution of a disoriented chiral condensate in 1+3 dimensions is not devoid of conceptual challenges.

- Whatever scenario one is willing to adopt, one should remember that the chiral condensate is an open quantum system. The implications of the existence of the corresponding hierarchy of time scales are worth being examined.

- The existence of Centauro-like events, if confirmed, would not be a mere curiosity. The creation and evolution of a disoriented chiral condensate is sensitive to the physics governing the early stage of the collision process and the experimental signature of the phenomenon is particularly unambiguous (the neutral/charged ratio). Although these events are presumably rare, they might convey relatively clean information.

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