Signals of Disoriented Chiral Condensate

Zheng Huang and Mahiko Suzuki
Theoretical Physics Group, Lawrence Berkeley Laboratory
University of California, Berkeley, California 94720, USA

Xin-Nian Wang
Nuclear Science Division, Lawrence Berkeley Laboratory
University of California, Berkeley, California 94720, USA

Abstract

If a disoriented chiral condensate is created over an extended space-time region following a rapid cooling in hadronic or nuclear collisions, the misalignment of the condensate with the electroweak symmetry breaking can generate observable effects in the processes which involve both strong and electromagnetic interactions. We point out the relevance of the dilepton decay of light vector mesons as a signal for formation of the disoriented condensate. We predict that the decay $\rho^0 \to \ell^+\ell^-$ will be suppressed and/or the $\rho$ resonance peak widens, while the decay $\omega \to \ell^+\ell^-$ will not be affected by the condensate.

PACS numbers: 11.30.Rd, 13.40.Hq, 13.85.-t
1 Introduction

Recently a proposal to observe the disoriented chiral condensate (DCC) in high energy collisions \[1, 2, 3, 4\] has received much attention. Both numerical and analytical solutions that may describe the formation of the DCC have been studied \[5, 6, 7\]. The basic idea of the DCC is that the collision debris will expand outward from the center of a collision at the speed of light, leaving behind a relatively cold interior whose vacuum has a chiral orientation different from that of the normal vacuum in the exterior. Unlike the quark-gluon plasma (QGP) which is usually assumed to be in thermal equilibrium, the DCC is not expected to be in a local thermal equilibrium. Indeed, the opposite extreme of free streaming of the collision debris is expected to produce a colder interior which leads to the DCC formation \[4\]. Rajagopal and Wilczek \[5\] suggested a quenching mechanism for formation of the DCC in which the temperature drops suddenly due to expansion. It thus seems that a zero temperature dynamics may be more appropriate to describe hadron dynamics in the DCC.

Perhaps the most direct test for an ideal DCC will be the observation of an anomalous isospin charge distribution \[1, 2, 3, 4\]

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

where \( r = n_{\pi^0}/(n_{\pi^+}+n_{\pi^-}) \). This prediction results from a long-range correlation in the isospin direction for a given event. It has been demonstrated numerically \[3, 7\] that a quench after collisions may actually yield a distinct cluster structure in the rapidity distribution which indicates a strong correlation in the isospin direction. In order for the DCC to be interesting, its size must be much larger than an inverse pion mass. But we do not know at present how long the correlation length will be for a typical DCC domain formed in hadronic or nuclear collisions. If the correlated domain is not large enough as compared to the total interaction region, especially in the central region of heavy ion collisions, the expected cluster structure in momentum phase space could be smeared out by the fluctuations of many domains. The distribution of the total neutral to charged pion ratio would have a binomial behavior rather than that of Eq. (1). It is thus important to study the consequences of DCC with both large
and small correlated domains.

In this paper we search for a possible signature of formation of the DCC, assuming that its space-time spread grows much larger than $1/m_\pi$. We will focus on the misalignment of the DCC with the direction of the electroweak symmetry breaking. When the hadronic resonances decay inside the DCC, their electromagnetic decay modes show a clear signal of the disorientation of the background. In contrast, purely hadronic decay modes are not affected by the disorientation in the limit of a small explicit chiral symmetry breaking and a slow space-time variation of the DCC. We shall study the dilepton decays of the $\rho$ and the $\omega$, since the DCC affects the $\rho$ and $\omega$ mesons quite differently because of their different chiral properties. In the limit of a slowly varying DCC, the $\rho$ resonance peak in the dilepton mass spectrum will be reduced by half. When the DCC varies rapidly in space-time, the $\rho$ resonance peak will be smeared and may even disappear. On the other hand, the $\omega$ resonance will be little affected by the DCC. Therefore there is a chance to test the formation of the DCC by carefully measuring the dilepton decays of the $\rho$ and the $\omega$ meson. We shall first study the limit of a slowly varying DCC, then include the space-time variation of the condensation.

2 Vacuum Misalignment

The DCC is created in a certain spatial domain, expands outward and eventually disappears. The vacuum orientation of the DCC labeled by an isovector rotation angle $\theta$ is in general space-time dependent, satisfying a boundary condition: $\theta = 0$ for $|x| > t$ (space-like) and $\theta \to 0$ for $t \to \infty$. Since such a configuration always has a higher energy than the true vacuum, the DCC vacuum decays into the true vacuum semiclassically by a coherent emission of pions, leading to the characteristic isospin charge distribution of (1). It is our hope that the DCC will live much longer than the typical time scale of low-energy strong interactions, and the DCC spreads over a spatial region much larger than $1/m_\pi$. Given a lifetime and a spatial size, the condensate may also vary with space-time, which determines the correlation of the emitted pions. An idealized limit often discussed in literature is an infinitely
large DCC which corresponds to an approximately homogeneous vacuum in the cool interior. We shall first study the DCC vacuum in this limit where the orientation angle \( \theta \) is space-time independent. We later extend our argument to the DCC of a finite size and also with a space-time variation in the interior.

Let us consider the DCC in the case of two flavors \((u \text{ and } d)\) in which the quark condensate matrices take the form

\[
\langle \bar{q}_L q_R \rangle = -v^3 \exp[i2\theta \cdot \tau]; \quad \langle \bar{q}_R q_L \rangle = -v^3 \exp[-i2\theta \cdot \tau].
\]  

(2)

The DCC vacuum is related to the normal one through a chiral rotation \(U_L\) and \(U_R\) with \(U_L^\dagger U_R = \exp[i2\theta \cdot \tau]\). It transforms the quark fields as

\[
q_L \rightarrow U_L q_L; \quad q_R \rightarrow U_R q_R.
\]

(3)

If the quark mass is ignored, the QCD vacua would be infinitely degenerate with respect to the chiral rotation. In this limit all strong interactions would take exactly the same form whichever vacuum is realized. When the quark mass is included, the degeneracy is lifted and the DCC vacuum energy density is

\[
E^{\text{QCD}}(\theta) = \langle \text{Tr}(\bar{q}_L U_L^\dagger M U_R q_R) \rangle + \text{c.c.}
\]

\[= -m_n^2 F_\pi^2 \cos 2\theta,
\]

(4)

where \(\theta = |\theta|\), and \(M\) is the diagonal mass matrix with eigenvalues \(m_u\) and \(m_d\). \(E^{\text{QCD}}(\theta)\) is minimized at \(\theta = 0\), that is, in the \(\sigma\) direction. However, this energy shift does not play an important role in the qualitative picture of the DCC. The strong interaction processes alone cannot tell of the chiral orientation of the condensate in the limit of the small current quark mass.

On the other hand, the misalignment of the DCC vacuum with the \(\sigma\) direction will be clearly seen through the electroweak interactions of hadrons. In the DCC where the quark fields are rotated as in (3), the electromagnetic current of the \(u\) and \(d\) quark is represented in the form

\[
J_{\mu}^{\text{EM}} \rightarrow \tilde{J}_{\mu}^{\text{EM}} = \bar{q}_L \gamma_\mu U_L^\dagger Q U_L q_L + \bar{q}_R \gamma_\mu U_R^\dagger Q U_R q_R,
\]

(5)
where $Q$ is the diagonal charge matrix with eigenvalues $2/3$ and $-1/3$, and the quark fields are in the basis where $\langle \bar{q}_L q_R \rangle_{DCC} = \langle \bar{q}_R q_L \rangle_{DCC} = -v^3$, i.e., in a basis in which the standard hadron mass spectroscopy is valid. The disorientation also raises the vacuum energy of the EM interaction, which is obtained by taking an expectation value of the effective Hamiltonian

$$\langle H \rangle_{DCC} = -\frac{e^2}{2} \int d^4x D^{\mu\nu}(x) \langle T \hat{j}_\mu^{EM}(x) \hat{j}_\nu^{EM}(0) \rangle,$$

where $D^{\mu\nu}$ is the photon propagator. When the explicit chiral symmetry breaking is ignored, the electromagnetic vacuum energy density is calculated

$$E^{EM}(\theta) - E^{EM}(0) = F_\pi^2 \frac{n_i^2 + n_2^2}{2} \sin^2 2\theta (m_{\pi^+}^2 - m_{\pi^0}^2),$$

where $n_i$ ($i = 1, 2, 3$) are the components of the unit vector $\mathbf{n} = \theta/\theta$. The true vacuum at $\theta = 0$ simultaneously minimizes both $E^{QCD}(\theta)$ and $E^{EM}(\theta)$. Therefore the normal QCD vacuum aligns with the direction of the spontaneous symmetry breaking of the electroweak interactions, which is fixed by the vacuum expectation value of the Higgs field. However, the DCC vacuum does not align with the electroweak vacuum. One immediate consequence of this misalignment is that the vector-meson dominance no longer holds for the EM current in the familiar way. In the normal vacuum, the isovector part of the EM current is dominated by the $\rho^0$ meson (770 MeV). In the DCC, the isovector current will be dominated by the mesons that would dominate the vector and axial-vector isospin currents in the normal vacuum. They are the $\rho$ meson and the $a_1$ meson (1260 MeV). The isovector EM current in DCC will show two resonance peaks, one at 770 MeV and the other at 1260 MeV. On the other hand, the isoscalar current is dominated by the $\omega$ meson (780 MeV) and the $\phi$ meson (1020 MeV) in the normal vacuum, which are the two singlets $(1, 1)$ of $SU(2)_L \times SU(2)_R$. Therefore, no matter how much the DCC vacuum is chirally disoriented, the $\omega$ and $\phi$ mesons are not affected at all. This difference between the isovector and isoscalar currents will result in very different behaviors of the dilepton decays of the $\rho$, $\omega$ and $\phi$ mesons.


3  Dilepton Decays of $\rho$ and $\omega$

In the normal vacuum, the dilepton decays of $\rho^0$ and $\omega$ occur through one photon with the effective couplings

$$\mathcal{L}_{\gamma V} = \frac{e}{2g_\rho} F_\mu^{(\gamma)} F^{(\rho)\mu}, + \frac{e}{6g_\omega} F_\mu^{(\gamma)} F^{(\omega)\mu},$$

where $F_\mu^{(\gamma)}$ is the field strength. On the mass-shell of $\rho^0$ and $\omega$, we can rewrite (8) into the effective interactions

$$\mathcal{L}_{\gamma V} = \frac{e m_\rho^2}{g_\rho} \rho^0 \mu A^\mu + \frac{e m_\omega^2}{3g_\omega} \omega \mu A^\mu.$$  

In the DCC, the meson state with which the isovector current couples is no longer a single meson state of mass 770 MeV but a linear combination of the states with mass 770 MeV and 1260 MeV ($= m_{a_1}$). When the $\rho$ and $a_1$ mesons dominate the vector and axial-vector isospin currents $J_\mu$ and $J_5\mu$, respectively, it holds in the normal vacuum [9] that

$$\langle 0 | J_5\mu | a_1 \rangle = \langle 0 | J_\mu | \rho \rangle = \frac{m_\rho^2}{g_\rho} \epsilon_\mu.$$  

Then the effective photon-meson couplings through the isovector current in the DCC are given by

$$\mathcal{L}_{\gamma V}^{DCC} = \frac{e m_\rho^2}{g_\rho} Tr \left\{ \frac{\tau_3}{2} \left[ U_L^\dagger (\rho_\mu - a_\mu) \cdot \tau U_L + U_R^\dagger (\rho_\mu + a_\mu) \cdot \tau U_R \right] \right\} A^\mu + \frac{e m_\omega^2}{3g_\omega} \omega \mu A^\mu,$$

The $\rho_\mu$ and the $a_\mu$ in (11) are the mass eigenstates of $\sim 770$ MeV and $\sim 1260$ MeV in the DCC, respectively, which are parity mixture states if a “parity” is defined in the normal vacuum. The same result can be obtained by assuming explicitly that the $\rho$ and $a_1$ fields transform as $\rho \pm a \sim (3, 1) \pm (1, 3)$ under the chiral rotation. To derive Eq. (11), one only needs a chiral transformation property of the isovector EM current, and the consequence (11) of a vector- and axial-vector-meson dominance for the chiral SU(2) currents. The latter requires that the $\rho$ and $a_1$ fields have the same
transformation properties as the currents, i.e., as \((3, 1) \pm (1, 3)\), leading effectively to a field-current identity \([10]\).

All possible charge states of the \(\rho\) and \(a_1\) mesons couple with the photon in the DCC with a general orientation. A violation of the charge conservation in these transitions is superficial. Since the DCC vacuum may be charged, an apparent violation of charge conservation in a photon-meson transition is compensated by a charged decay of the vacuum that follows. It is reasonable to expect that when the DCC is created, all possible orientations of the chiral SU(2)\(_L\)×SU(2)\(_R\) are allowed with an equal probability. We average the photon-meson transitions in the group-parameter space of SU(2)\(_L\)×SU(2)\(_R\). On average, all six transitions, \(\rho \to \gamma\) and \(a_1 \to \gamma\), occur with an equal probability;

\[
\langle |C^i_{\rho}(\theta_L, \theta_R)|^2 \rangle = \langle |C^j_{a_1}(\theta_L, \theta_R)|^2 \rangle = \frac{1}{6}, \quad (i, j = 1, 2, 3) \tag{12}
\]

This means that the \(\rho\) peak in the \(\ell^+\ell^-\) mass plot will be reduced by half if all three charge states of \(\rho\) are produced with an equal rate, and a new broad bump of width \(\simeq 300\) MeV (= \(\Gamma_{a_1}\)) may appear at 1260 MeV

\[
\Gamma (\rho(770\text{MeV}) \to \ell^+\ell^-)_{\text{DCC}} = \frac{1}{2} \Gamma (\rho^0 \to \ell^+\ell^-)_{\text{normal}}, \tag{13}
\]

\[
\Gamma (a_1(1260\text{MeV}) \to \ell^+\ell^-)_{\text{DCC}} = \left(\frac{m_{\rho}}{m_{a_1}}\right)^3 \Gamma (\rho(770\text{MeV}) \to \ell^+\ell^-)_{\text{DCC}} \tag{14}
\]

This prediction is little affected by the current quark mass. It is a consequence of the misalignment of the DCC with the electroweak vacuum fixed by the Higgs field, so it would disappear only at \(F_\pi \to 0\) in which a distinction between the \(\rho\) and the \(a_1\) ceases to be meaningful.

In a real experimental environment, the reduction of the \(\rho\) peak will be lessened by the probability of the DCC formation in collisions. Note, however, that the averaging in Eq. (12) is done over all possible states including the normal vacuum as a special case and those vacua that are only slightly different from it. If we can select the likely candidates for the DCC events by the isospin charge asymmetry \([1, 2, 3, 4]\), the reduction of the \(\rho\) peak should be even larger. However, if the \(\rho\) and \(\omega\) mesons are copiously produced together with the coherent pions emission at the time of the decay of the DCC itself, the test may become difficult.
The above analysis has so far been based on the idealized limit of an infinitely large and slowly varying DCC. Though very little is known about the detailed properties of the DCC, we should examine how the argument is modified by the space-time variation of the DCC.

4 Space-time Variation Effects

When the size of the DCC is finite in space-time, it can happen that the vacuum acts like a reservoir of energy and momentum for hadrons and leads to an apparent violation of energy and momentum conservation in their decay and scattering processes. The energy and momentum transfer may become larger when the DCC also has a space-time variation in the interior. For a “realistic” DCC, we must take into account these space-time variations of the DCC. The magnitude of the transfer, when it occurs, is determined by the uncertainty principle:

\[ \Delta p \sim \frac{1}{\Delta x} ; \quad \Delta E \sim \frac{1}{\Delta t} , \]  

(15)

where \( \Delta x \) and \( \Delta t \) are the space and time variations. When the DCC of a finite size is slowly varying in the interior, \( \Delta x \) and \( \Delta t \) will be the space-time size of the DCC. For the DCC with a rapid space-time variation in the interior, \( \Delta x \) and \( \Delta t \) are determined by the scale of the variation. A DCC is meaningful and interesting only when its size is much larger than \( 1/m_\pi \), while the space-time variation in the interior may be as large as \( 1/m_\pi \). Let us parametrize \( \Delta x \) and \( \Delta t \) as

\[ \Delta x \simeq \Delta t \equiv \frac{\kappa}{m_\pi} . \]  

(16)

This amounts to a violation of energy-momentum conservation with a magnitude \( \Delta p \simeq \Delta E \leq m_\pi / \kappa \) in decay processes and in final-state interactions, if they exchange energy-momentum with the vacuum. The value of \( \kappa \) ranges from \( \text{O}(1) \) to \( \sim 5 \) or more, depending on whether it is determined by the variation in the interior or by the size of the DCC. We will examine how this space-time dependence changes the previous oversimplified picture.
To be concrete, we describe $\rho$, $a_1$, and $\omega$ as the massive gauge-like bosons of $SU(2)_L \times SU(2)_R$ and of $U(1)_V$. It has been long recognized that this is phenomenologically a good description of the spin-one mesons \cite{10}. The relevant Lagrangian is that of a gauged $\sigma$ model to which the vector meson masses are added:

$$
\mathcal{L} = \text{Tr}D_\mu \Phi \dagger D_\mu \Phi - \frac{1}{4} \text{Tr} (F_{L\mu\nu} F_{L}\^{\mu\nu} + F_{R\mu\nu} F_{R}\^{\mu\nu}) - \frac{1}{4} \omega_\mu \omega_\mu \\
+ \frac{m_\rho^2}{2} \text{Tr} (V_{L\mu} V_{L}^{\mu} + V_{R\mu} V_{R}^{\mu}) + \frac{m_\omega^2}{2} \omega_\mu \omega_\mu - U_\omega - U(\Phi \dagger \Phi) \quad (17)
$$

where $\Phi = \frac{1}{2} (\sigma + i \pi \cdot \tau)$, $V_{L(R)\mu} = \frac{1}{2} (\rho_\mu \pm a_\mu) \cdot \tau$ and

$$
D_\mu \Phi = \partial_\mu \Phi + igV_{L\mu} \Phi - ig \Phi V_{R\mu} . \quad (18)
$$

In (17), $U_\omega$ is the anomaly term that generates the $\omega \to \pi \pi$ decay amplitude, and $U(\Phi \dagger \Phi)$ denotes the mass and potential terms of the $\Phi$. Upon spontaneous symmetry breaking, the $a_1$ meson acquires an additional mass. The explicit symmetry breaking term $-m_\pi^2 F_\pi \sigma$ is left out for the moment. We will discuss its effect later.

For the DCC with the space-time dependent angle parameters $\theta_L(x)$ and $\theta_R(x)$, the quarks fields $q'(x)$ in the DCC are related to the quark fields $q(x)$ in the normal vacuum through

$$
q_L = U_L(x)q'_L ; \quad q_R = U_R(x)q'_R , \quad (19)
$$

where $U_L(x)$ and $U_R(x)$ depend on space and time instead of global transformations in Eq. (3). The same transformation relates the meson field $\Phi'$ in the DCC to the meson field $\Phi$ in the normal vacuum as

$$
\Phi = U_L(x)\Phi' U_R^{\dagger}(x) ; \quad \Phi^{\dagger} = U_R(x)\Phi'^{\dagger} U_L^{\dagger}(x) , \quad (20)
$$

where $\langle \Phi \rangle = F_\pi/2$ in the normal vacuum. In a space-time dependent DCC, the spin-one fields also develop vacuum expectation values since the background carries angular momentum. One convenient basis for the spin-one fields is obtained by the local transformations,

$$
V_{L\mu} = U_L(x)V_{L\mu} U_R^{\dagger}(x) - \frac{i}{g} U_L(x)\partial_\mu U_R^{\dagger}(x) \\
V_{R\mu} = U_R(x)V_{R\mu} U_L^{\dagger}(x) - \frac{i}{g} U_R(x)\partial_\mu U_L^{\dagger}(x) , \\
\omega_\mu = \omega'_\mu \quad (21)
$$
where $\langle V_{L(R)\mu} \rangle = 0$ in the normal vacuum. Since (21) is an SU(2)$_L \times$SU(2)$_R$ gauge transformation, the Lagrangian in (17) is invariant except for the $\rho$-$a_1$ mass term

$$L \rightarrow L + \delta L,$$

$$\delta L = i \frac{m^2_\rho}{g} \text{Tr} \left[ U_L^\dagger(x) \partial^\mu U_L(x) V_{L\mu}' \right] + i \frac{m^2_\rho}{g} \text{Tr} \left[ U_R^\dagger(x) \partial^\mu U_R(x) V_{R\mu}' \right] + \frac{m^2_\rho}{2g^2} \text{Tr} \left[ \partial^\mu U_L(x) \partial_\mu U_L^\dagger(x) \right] + \frac{m^2_\rho}{2g^2} \text{Tr} \left[ \partial^\mu U_R(x) \partial_\mu U_R^\dagger(x) \right].$$

In this basis, $\delta L$ contains a space-time dependent tadpole interaction for $\rho$ and $a_1$. The presence of these tadpoles means that $\rho$ and $a_1$ acquire effectively space-time dependent mass in the DCC. One way to understand this interpretation is to sum an infinite series of the tadpole diagrams in the $\rho$ and $a_1$ propagators. Alternatively we can eliminate the tadpole terms by using the exact classical equations of motion for the $\rho$ and $a_1$ fields in the presence of the $\Phi$ field. This amounts to an additional shift $\langle V_{L(R)\mu}' \rangle$ for the $\rho$ and $a_1$ fields in addition to (21), which generates additional vector-meson mass terms through the interaction $g^2 V_{\mu} V_{\nu} V_{\mu} V_{\nu}$, contained in $\text{Tr}(F_{\mu\nu} F^{\mu\nu})$. We can see from $\delta L$ that the DCC expectation value of $V_{L(R)\mu}'$ is related to the space-time variation of the DCC

$$\langle V_{L(R)\mu}' \rangle_{\text{DCC}} = O \left( \langle U_L^\dagger(x) \partial_\mu U_L(x) \rangle_{\text{DCC}} \right) \simeq \frac{m_\pi}{\kappa}.$$ (23)

This additional shift generates also a new $\rho \rightarrow \pi\pi$ interaction with the coupling of $O(g^2 m_\pi/\kappa m_\rho)$ through $g^2 (\langle \rho_\mu \rangle \times \pi) \cdot (\rho^\mu \times \pi)$. This is however very small when $\kappa \geq 1$. The space-time dependent part of the mass is $\Delta m_\rho^2 \sim g^2 \langle V_{\mu} \rangle \langle V^\mu \rangle$. With $g^2/4\pi = 2.5$ from the $\rho \rightarrow \pi\pi$ decay width, the space-time dependent smearing of the $\rho$ mass is of the order of $\Delta m_\rho^2 \sim 30m_\pi^2/\kappa^2$. The smearing in mass square translates to a smearing in mass itself as

$$\delta m_\rho \leq \frac{30m_\pi^2}{2\kappa^2 m_\rho}.$$ (24)

If the DCC is slowly varying and the size of the DCC determines $\kappa$ ($\geq 5$), $\delta m_\rho \leq 25$ MeV; the $\rho$ meson will maintain its resonance shape with a suppressed $\ell^+\ell^-$ event number. When the space-time variation in the interior is rapid ($\sim O(1/m_\pi)$), the $\rho$
resonance becomes very wide due to the smearing effect, and may virtually disappear if most of $\rho$'s are created in contact with the vacuum. The effective $\rho$-$\gamma$ transition coupling also acquires a space-time dependence, which adds to an apparent violation of energy-momentum conservation in the decay $\rho \rightarrow \ell^+ \ell^-$. Therefore, in the case of a rapidly varying DCC, one may expect a rather prominent shape change at the $\rho$ resonance in both the $\pi\pi$ decay mode and the $\ell^+ \ell^-$ decay mode. In both cases, the $\omega$ field does not change and acts as if there were no DCC background. In particular, the $\omega$ resonance peak of the dilepton will maintain its shape and serves as a “standard”. A careful comparison of the $\rho$ and the $\omega$ components at the overlapped $\rho$-$\omega$ peak will reveal the existence of DCC.

We finally remark on the explicit chiral symmetry breaking that has so far been ignored. The main effect of this breaking is to smear the effective pion mass square over a range from $-m_\pi^2$ to $m_\pi^2$ in the DCC. Due to its short lifetime, the $\rho$ decays mostly inside the DCC. The smearing of the pion mass has very little effect on the $\rho$ decay width because the phase space for $2\pi$ is large enough. The effect on the phase space for $3\pi$ in the decay $\omega \rightarrow \pi\pi\pi$ will be a little larger. If the decay $\omega \rightarrow \pi\pi\pi$ occurs inside the DCC, some widening of the $\omega$ width may occur. However, the $\omega$ lifetime is probably too long for the $\omega$ to decay inside the DCC. If the majority of the $\omega$ decays outside the DCC, the $3\pi$ peak of the $\omega$ will remain relatively clean.

5 Discussions

We have shown that the misalignment of the DCC with the electroweak direction results in the suppression of the decay $\rho^0 \rightarrow \ell^+ \ell^-$ by a factor of two on average, while no suppression is expected for $\omega \rightarrow \ell^+ \ell^-$. The marked difference between $\rho$ and $\omega$ arises from their different chiral properties. This is in sharp contrast to the QGP, which affects equally the $\rho^0 = \frac{1}{\sqrt{2}}(\bar{u}u - \bar{d}d)$ and the $\omega = \frac{1}{\sqrt{2}}(\bar{u}u + \bar{d}d)$ except for the large difference in their lifetimes which favors $\rho^0 \rightarrow \ell^+ \ell^-$ over $\omega \rightarrow \ell^+ \ell^-$ during the hadronic expansion [11]. Some may question whether the $\rho$ and the $\omega$ can be produced at all in the QGP. It appears that the DCC is a less hostile environment for these resonance to be formed. In any case, we have to learn about it from experiment. It is clearly
necessary to carry out more detailed theoretical analysis of the outgoing dilepton spectrum for the space-time dependent DCC of which we have little understanding at present.

In order to carry out our test in experiment, we need to know the relative production rate of $\rho$ and $\omega$ in hadron collisions. At very low energies, it is known [12] that the production rates for the $\rho^0$ and the $\omega$ are equal in pp collision: $\sigma(\rho^0)/\sigma(\omega) = 1.0 \pm 0.2$ at $\sqrt{s} = 5$ GeV; $\sigma(\rho^0)/\sigma(\omega) = 1.07 \pm 0.2$ at $\sqrt{s} = 6.8$ GeV. Recently, the cross sections were measured at higher energy by NA27 experiment [13] through the hadronic decay modes. The result is $\sigma(\rho^0) = (12.6 \pm 0.6)$ mb vs. $\sigma(\omega) = (12.8 \pm 0.8)$ mb in pp collision at $\sqrt{s} = 27.5$ GeV. The quark model predicts that these cross sections ought to be equal. Though an independent check for their separate cross sections would be desirable, the equality of the production rates appears to be a very safe assumption.

A search for the Centauro events has been made by UA5 [14] and UA1 [15] at $\sqrt{s} = 546$ GeV and by UA5 [16] at $\sqrt{s} = 900$ GeV. An upper limit was set on the Centauro production by UA5 [16] at the level of a few times $10^{-3}$ per inelastic events at $\sqrt{s} = 900$ GeV. Since an ideal DCC with large correlated domains will generate the Centauro and anti-Centauro events, there has been so far no evidence for clear DCC events. We hope that our suggestion made in this paper will be considered in future experiments searching for DCC, in both hadronic and heavy nucleus collisions.

Acknowledgements

We are grateful to J. D. Bjorken for stimulating conversations. This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Divisions of High Energy Physics and Nuclear Physics of the U.S. Department of Energy under contract DE-AC03-76SF00098 and in part by the U.S. National Science Foundation under grant PHY90-21139. Z.H. acknowledges a financial support of the National Science and Engineering Research Council of Canada.
References

[1] A.A. Anselm, Phys. Lett. B 217, 169 (1989); A.A. Anselm and M.G. Ryskin, Phys. Lett. B 266, 482 (1991).

[2] J.-P. Blaizot and A. Krazywicki, Phys. Rev. D 46, 246 (1992).

[3] J. Bjorken, Int. J. Mod. Phys. A 7, 4189 (1987); Acta Physics Polonica B 23, 561 (1992); J. Bjorken, K. Kowalski and C. Taylor, SLAC Preprint SLAC-PUB-6109 (1993).

[4] J. Bjorken, K. Kowalski and C. Taylor, SLAC Preprint, SLAC-PUB-6109 (1993).

[5] K. Rajagopal and F. Wilczek, Nucl. Phys. B 399, 395 (1993); B 404, 577 (1993); F. Wilczek, Princeton Preprint IASSNS-HEP-93/48 (1993).

[6] S. Gavin, A. Gocksch and R.D. Pisarski, Brookhaven Preprint BNL-GGP-1.

[7] Z. Huang and X.N. Wang, LBL Preprint LBL-34931 (1993).

[8] M.E. Peskin, Nucl. Phys. B 185, 197 (1981); J. Preskill, Nucl. Phys. B 177, 21 (1981).

[9] S. Weinberg, Phys. Rev. Lett. 18, 507 (1967)

[10] N. Kroll, T.D. Lee and B. Zumino, Phys. Rev. 157, 1376 (1967); T.D. Lee and B. Zumino, Phys. Rev. 163, 1667 (1967).

[11] U. Heinz and K.S. Lee, Nucl. Phys. A 544, 503c (1992).

[12] V. Blobel et al., Phys. Lett. B 48, 73 (1974); Nucl. Phys. B 69, 237 (1974).

[13] M. Aguilar-Benitez et al., LEBE-EHS Collaboration, Z. Phys. C 50, 405 (1991)

[14] K. Alpgard et al., UA5 Collaboration, Phys. Lett. B 115, 71 (1982).

[15] G. Arnison et al., UA1 Collaboration, Phys. Lett. B 122, 189 (1983).

[16] G.J. Alner et al., Phys. Lett. B 180, 415 (1986).