HOW BLACK IS A CONSTITUENT QUARK??

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Dedicated to Wiesław Czyż in honour of his 65th birthday

We review the question posed in the title, with emphasis on hard-collision phenomena utilizing rapidity gaps and jets as a diagnostic. If constituent quarks are very black, it is conceivable that various exotic phenomena exist, including production of strange matter, disoriented chiral condensate, and/or ultra-hot quark-gluon plasma.

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1. Introduction

The title and much of the subject matter of this paper is not very new. The question was debated in the late 1970s and early 1980s after the advent of the additive quark [1], especially in Poland [2] and in St. Petersburg [3]. Nevertheless the issue seems to remain unresolved and becomes of more immediate importance as the energy scale of colliders increases [4]. In particular we may expect on very general grounds that the blackness of anything increases as its energy is increased. My favorite argument is that a reasonable measure of interaction probability is the momentum density in the impact-plane (in GeV/fm^2) of the Lorentz-contracted pancake representing the projectile. The dependence of interaction probability on momentum density is undoubtedly monotonically increasing (Fig. 1(a)). Then upon a Lorentz boost of one projectile by, say, a factor 10, we see that for any reasonable distribution of momentum density the central opacity, as well as the interaction radius, must increase.

At the SSC, collisions of single constituent quarks will be seen up to cms energies $\lesssim$ 20 TeV (corresponding to their $x_F \lesssim 0.5$). Even within

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the nominal short-distance domain of perturbative QCD, nonperturbative evolution of blackness is anticipated [4]. Therefore it seems especially timely to raise the issue again.

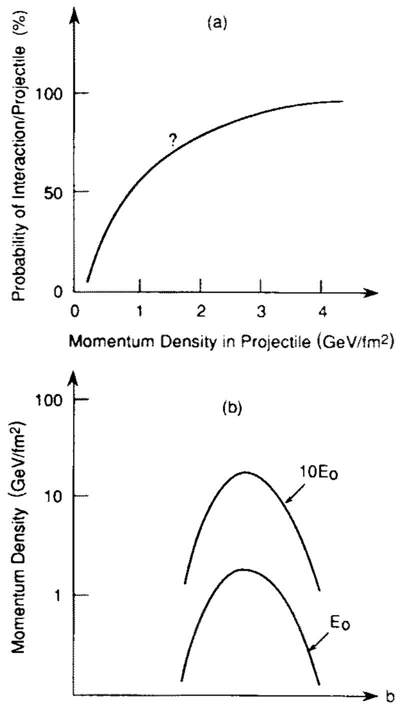

Fig. 1. Why blackness and interaction radii increase with energy. Note the interaction probability is the product of the probabilities (a) for each projectile; this in turn is invariant under longitudinal boosts.

We shall here simplify the situation by imagining that the pp collisions are replaced by heavy-quark BB collisions. The advantage is that there is only one constituent quark in each projectile and complications of rescattering, shadowing, three-body wave functions, etc., are absent, and the essential points emerge more clearly. The disadvantage is that there never will be any data. However, it is not difficult to imagine how data should look, so that general guidance to the real problem is still available. After a quick reminder of the Isgur–Wise heavy-quark limit of QCD [5], we apply it in Section 2 to elastic and diffraction scattering of B's to see what if anything is learned about the internal structure of the quark. We shall find a somewhat disappointing result, which impels us onward in Section 3 to the consideration of hard-collision processes.

In Section 4 we briefly review the expectations from perturbative QCD for local blackness around leading partons (the "small-x" problem [4], described by the BFKL equation). Finally in Section 5 we speculate on what might happen at extreme energies when extreme-relativistic opaque matter comes into collision. The possibilities are quite speculative, but include disoriented chiral condensate, strange-matter nuggets, and quark-gluon plasma with extraordinarily high initial temperature.

2. The heavy quark limit of QCD and B-B scattering

We view [5] a B meson as a constituent quark of rather small size (0.2–0.3 f) orbiting a very heavy b-quark, whose interactions are unimportant — other than providing a binding potential for the light quark [6]. We shall ignore spin. Now consider B-B scattering in some center-of-mass frame. The velocities of the heavy b quarks (and their mesons) are unaffected by the collision. The scattering amplitude in momentum space in general has the form (Fig. 2) (neglecting[1] spin!)

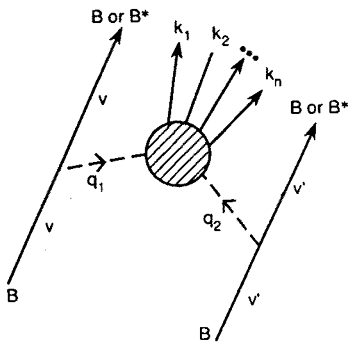

Fig. 2. Kinematics for B-B collisions; it is really *by definition* the collision of the light constituent quarks.

$$T = T(v, v'; k_1, \ldots, k_n; q, q'), \tag{1}$$

where $\{k_i\}$ are the four-momenta of the produced particles and q and q' are the momenta transferred by the heavy quarks; they satisfy

$$q \cdot v = q' \cdot v' = 0. \tag{2}$$

The phase space sum is

$$\frac{(2\pi)^4}{(2\pi)^{3(n+2)}} \frac{d^3k_1}{2\omega_1} \cdots \frac{d^3k_n}{2\omega_n} d^4q\, d^4q'\, \delta(q\cdot v)\, \delta(q'\cdot v')\, \delta^4(q+q'-\Sigma k_i)\,. \tag{3}$$

[1] While we cheat here, the considerations are correct for Λ_b–Λ_b scattering.

For elastic scattering all of that boils down to $d^2q_\perp$, as expected. In that case (as perhaps in others) it is best to Fourier transform to impact-space. A naive calculation of the amplitude is then simply given by a convolution of wave-functions (Fig. 3)

$$T(B) = \int d^2b_1\ d^2b_2\, d^2b\ \rho(b_1)\ \rho(b_2)\ t(b)\ \delta^2(B + b_2 + b - b_1)\,, \tag{4}$$

where $\rho(b)$ is the square of the constituent quark wave function about the b quark, normalized such that

$$\int d^2b'\,\rho(b') = 1\,. \tag{5}$$

To get the cross section

$$(\text{const})\cdot \text{Im}\, T(B) = \frac{d\sigma_{\text{BB}}}{d^2B} = \int d^2b\, d^2b_1\, \rho(b_1)\, \rho(b_1 - B)\, \frac{d\sigma_{\text{qq}}}{d^2b} \cong \sigma_{\text{qq}}\rho_{\text{BB}}(B) \tag{6}$$

with

$$\int \rho_{\text{BB}}(B) d^2B = 1\,. \tag{7}$$

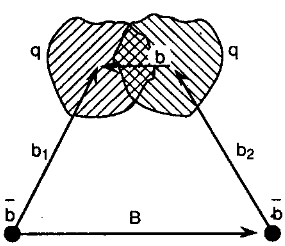

Fig. 3. Collision geometry in the impact plane for B-B scattering.

We have assumed the interaction range between the quarks is small compared to the meson size. Therefore the details of the impact-parameter dependence do not show up in σ_{el} or σ_{tot}. This is not at all surprising. To learn about the blackness of the quarks by the usual methodology we would have to see a shadow cast by them in their elastic scattering—after somehow making a coarse-grained average over the excited "atomic" states of the B mesons and probably over t as well. But even defining what is meant by the "elastic" qq final states is fraught with ambiguity. By the time t is large enough to be useful, relativistic kinematics and multiparticle production also becloud the issue.

These complications already occur for B-B scattering. Proton-proton scattering will be still more complicated. One might hope that those complications could be turned to one's advantage. That has been tried [2], but lots of ambiguity persists. We therefore pursue this line no further and turn to hard collisions.

3. Hard collisions and black quarks

Most hard collisions in B-B interactions will be dominated by those initial-state configurations in which the quarks themselves collide. If we sum over impact parameter B between the b-quarks, the result is simply

$$\frac{d\sigma_{ab}}{d^2b} \sim \sigma_{\text{Hard}} \int d^2b'\rho_a(b')\,\rho_b(b-b') \equiv F_{ab}(b)\,\sigma_{\text{Hard}}\,, \tag{8}$$

where ρ_a and ρ_b are the densities of the partons a and b which initiate the hard collision. These partons should reside *within* the constituent quarks;[2] we assume that the momentum transfer Q^2 is so large that the impact-parameter separation of these initiating partons is small compared to the constituent quark size ($Q^2 \gg 1$ GeV^2).

Given a hard process such as this as a "tag" of a close encounter of the constituent quarks, one may now search for an additional phenomenon in the same collision which can provide a measure of blackness. One possibility is a second hard collision in the same event. If we (1) let the ρ's in the above equation denote the densities of partons with x_i large enough to initiate the hard process of interest, and (2) assume that they form an uncorrelated distribution within the constituent quark, it follows that [7]

$$\frac{d\sigma^{\text{Double}}}{d^2b} = \frac{F_{ab}^2(b)}{2}\,\sigma_{\text{Hard}}^2 \tag{9}$$

and

$$\sigma_{\text{Double}} = \frac{\sigma_{\text{Hard}}^2}{2\sigma} \tag{10}$$

with

$$\frac{1}{\sigma} = \frac{\int d^2b\,[F_{ab}(b)]^2}{[\int d^2b\,F_{ab}(b)]^2}\,. \tag{11}$$

With Gaussian distributions for the convolution

$$F(b) = e^{-b^2/R^2} \tag{12}$$

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[2] If their longitudinal fractions are sufficiently large, this is a near certainty.

one gets

$$\frac{1}{\sigma} = \frac{1}{2\pi R^2}, \tag{13}$$

namely a number of order the quark-quark cross-section.

It is not clear what is learned from this exercise. First of all, hadron collider data seem not to show a very large double-parton cross-section; UA2 [8] finds σ to be bounded from below by 8.3 mb (95% C.L.), although the AFS collaboration at the ISR [9] measures σ to be of order 5 mb. And the assumption of uncorrelated parton distributions is suspect, not so much with respect to impact parameter as to momentum fraction x. One must be at small enough x to be sure that presence of one large-x parton does not inhibit the existence of another. For example, it is impossible for two partons in the same constituent quark to simultaneously have $x > 0.5$. On the other hand, at very small x the quark radius is sure to grow.

And whatever the outcome, the measurement does not directly address the question of blackness, which has to do with the frequency of the very commonplace, soft interactions which accompany the hard process, not the relatively rare second hard process.

A better class of processes — perhaps the best — is hard double diffraction, ideally as created by electroweak boson exchange. This is the process that initiated this author's interest in the question [10]. The issues are discussed elsewhere [11] and we only summarize here the situation for B-B scattering. Consider the two-photon-exchange process shown in Fig. 4, with subprocess

$$q + q \to q + q + \mu^+\mu^- . \tag{14}$$

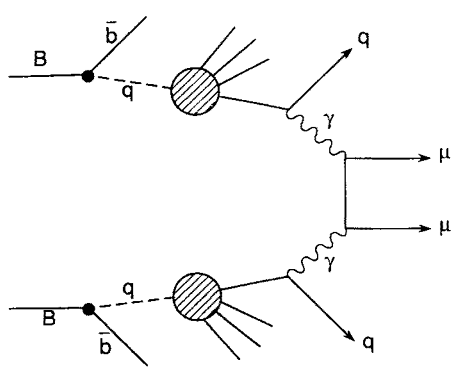

Fig. 4. The process of dilepton production into a rapidity-gap. The dashed lines describe the constituent quarks.

The quarks are partons within the constituent quarks, and we assume the photons are quite virtual: $q_1^2, q_2^2 \gg 1$ GeV^2. This leads to an event topology shown in Fig. 5. It is essentially the Drell–Yan dilepton signature, with the additional requirements that the dilepton be produced within a rapidity gap, and that tagging jets with $p_{t_1} \approx \sqrt{q_1^2}$; $p_{t_2} \approx \sqrt{q_2^2}$ appear at the edges of the rapidity gap.

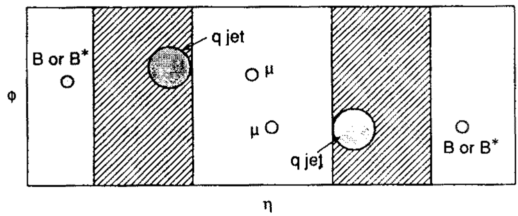

Fig. 5. Event structure in the lego plot for the process of Fig. 4.

The main point here is that the quark-partons collide head-on; hence the constituent quarks must also have a close encounter. However the presence of the rapidity-gap signals the *absence* of any additional spectator interactions and therefore measures the transmission probability of the quarks passing through each other. More precisely, we have for the probability of the hard collision (Fig. 6).

$$d\sigma_{\text{hard}} = \int d^2b_1 \; \rho_B(b_1) \; d^2b_2 \; \rho_B(b_2) \; d^2b_1' \; \rho_q(b_1') \; d^2b_2' \; \rho_q(b_2') \times \delta^2(B + b_2 + b_2' - b_1 - b_1') \, \sigma_{\text{hard}} \, . \tag{15}$$

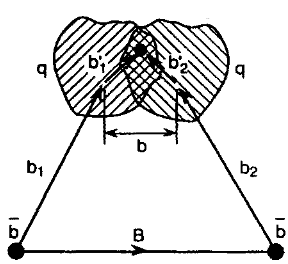

Fig. 6. Impact-plane geometry for the hard-collision process involving constituent quarks.

The cross-section including survival of the gap is obtained, in the absence of further correlation, by multiplying by two transmission probabilities. The first is simply $|S_{qq}(b)|^2$, the probability the two constituent quarks penetrate each other without (further) interactions. This is just what we need to estimate the blackness, as we shall show in a minute. The second factor is the same for the B-mesons, $|\tilde{S}(B)|^2$, with one caveat. In order to avoid double-counting, $|\tilde{S}(B)|^2$ should only account for that part of the B-B interaction *not* described by the additive quark-model, *e.g.*, the string-string interaction (Fig. 7). In an eikonal description these contributions are additive in the exponent; hence the factorized structure we assume would naturally come out of a theoretical analysis. But it may be more difficult to implement an experimental separation. In principle the supplementary interactions such as the string-string interaction could be calibrated from Υ-Υ collisions. But in any case the correction is probably not too big — and the sign is known. Ignoring it gives an *upper bound* on the blackness of the constituent quarks. The prediction is

$$\frac{\sigma_{\rm gap}({\rm BB} \to {\rm qq}\mu\mu X)}{\sigma_{\rm tot}({\rm BB} \to {\rm qq}\mu\mu X)} = \left\langle |\tilde{S}|^2 \right\rangle_{\rm BB} \cdot \left\langle |S|^2 \right\rangle_{\rm qq} \tag{16}$$

with

$$\left\langle |S|^2 \right\rangle_{\rm qq} \cong \frac{\int d^2b\, F_{ab} |S_{\rm qq}(b)|^2}{\int d^2b\, F_{ab}(b)} \,. \tag{17}$$

If we choose

$$|S|^2 = \exp\{-\nu\}\ \exp\{-b^2/R^2\}\,, \qquad F_{ab} = \exp\{-b^2/R^2\}\,, \tag{18}$$

the integrations give

$$\left\langle |S|^2 \right\rangle = \frac{1}{\nu}\left(1 - \exp\{-\nu\}\ \right) . \tag{19}$$

Even for quite large ν, there remains a considerable survival probability $\left\langle |S|^2 \right\rangle$ because almost all distributions of density have a lot of gray edge. More discussion of such convolutions can be found in Ref. [11].

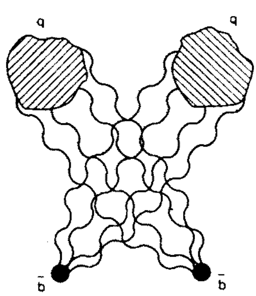

Fig. 7. Schematic picture of the "string-string" interaction.

Another example of a process sensitive to quark blackness is production of a Higgs boson within a rapidity gap [11]. However one would prefer to use that process for other purposes first!

4. Blackening of partons

Even in the regime of perturbative QCD there is a mechanism for blackening of partons as the energy increases. Indeed as $s \to \infty$ for fixed large t, say $t \sim 100\ \mathrm{GeV}^2$, it is expected [12] that the quark-quark interaction becomes *strong*. Mueller and Navelet [13] have described experiments to test this idea. They correspond to the separate measurement of total and elastic cross-sections. The first experiment is the measurement of

$$\frac{d\sigma}{d\eta_1 d\eta_2 d\ln Q_1^2 d\ln Q_2^2 d\phi_{12}} \tag{20}$$

for two jets of transverse momenta $Q_1^2 \sim Q_2^2 \sim t$ and originating from leading quarks, *i.e.*, at the extrema of the lego plot as allowed by kinematics. This is inelastic q-q scattering allowing for multiple production of jets (Fig. 8). The expectation, based on the properties of the BFKL evolution equation [14] for the multiple production of jets, is that the above cross-section, integrated over azimuthal angles, is enhanced by orders of magnitude over what is expected from single gluon exchange. A sketch of the behavior of the enhancement factor is shown in Fig. 9.

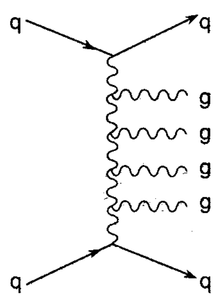

Fig. 8. Multijet production in q-q interactions.

The second Mueller–Navelet experiment is the analogue of elastic qq scattering. One searches for a coplanar component in the azimuthal distribution (Fig. 10) which corresponds to the subprocess in Fig. 11, with color singlet in the t-channel, and no jets in the middle of the lego plot. A fraction $\langle |S|^2 \rangle$ of these events in fact should have a rapidity gap. This fraction is a good signature and also might be yet another way of measuring the quark blackness.

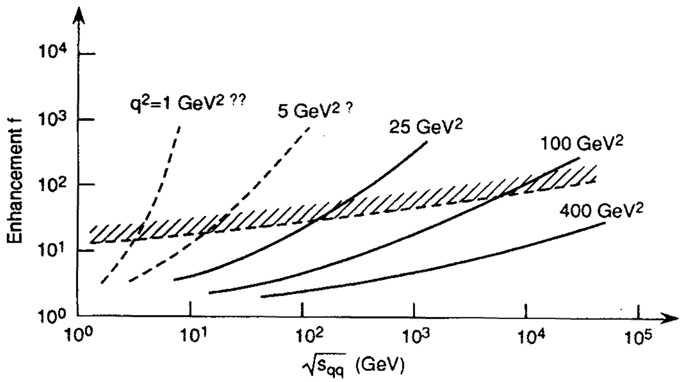

Fig. 9. My rough sketch of the enhancement f as described by Mueller and Navelet [14]. This is a free adaptation, for which those authors are not to blame. The region above the horizontal shaded band is subject to higher order corrections, not well understood theoretically.

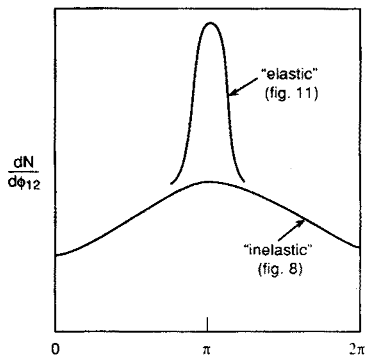

Fig. 10. Expected azimuthal correlation of quark jets in q-q scattering at very large s.

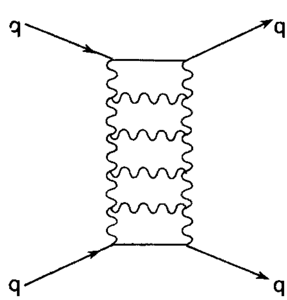

Fig. 11. The "elastic" contribution to Fig. 10. The fraction of area under the elastic peak is σ_{el}/σ_{tot}, not too small for black quarks.

The domain of legitimacy of the BFKL equation is such large values of s that $\alpha_s(t) \ln s \gg 1$. However, from inspection of Fig. 9 it is a little hard to see how or why the enhancement should not qualitatively look like the dashed contributions. In order words if one does have black constituent quarks when t is ~ 1 GeV^2 and s is ~ 100 GeV^2, it may simply be the anticipation of the phenomenon that perturbative QCD more reliably predicts at higher s and t.

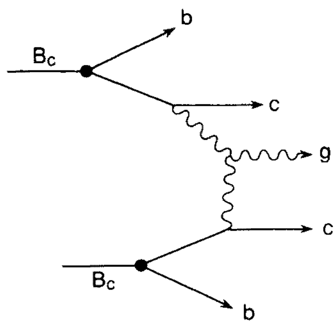

Fig. 12. B_c–B_c scattering as an easy way to enter the world of BFKL.

What is the physics? The equations are complicated [14] and at SLAC we struggle to attain a simple qualitative understanding. Consider B_c-B_c scattering at high energies. Beyond the one-gluon exchange there will be a radiative correction (Fig. 12). Viewed from the B_c rest-frame it is a Bethe–Heitler bremsstrahlung process. The gluon cloud is created by the dipole moment of the source. But after much cascading of this mechanism some soft gluons diffuse outside the radius b of interest leaving an exposed net color charge as well as dipoles in the core of the cascade. Thus this branching leads to a more intense core, which in turn leads to stronger gluon emission and escalation of the problem. In any case, the experimental study of this expected enhancement, which requires large qq subenergies in the TeV scale and above, will be of considerable interest.

5. What might black quarks do?

It is fun to speculate as to what might happen if quarks really do become black at SSC energies. If we interpret ν as the "number of mean free paths," or the "number of branchings" in some branching or cascading process, or the "number of wounds per parton," [15] then the energy-scale of survivor partons relative to the initial energy is $\sim 2^{-\nu}$. So we might expect strong absorption when

$$\frac{(1\ \mathrm{GeV})^2}{E_1 E_2} \lesssim (2^{-\nu_1})(2^{-\nu_2}) \sim \exp\{-(\nu_1+\nu_2)\ln 2\}. \qquad (21)$$

That is, if we ever have the condition

$$\nu = \nu_1 + \nu_2 \gtrsim \ln\left(\frac{s}{1\ \mathrm{GeV}^2}\right) \tag{22}$$

it becomes thinkable that cascading reaches "shower maximum" and a significant fraction of the incident energy-momentum is thermalized, or at the least arrested. Here we naively confine our attention to the beam fragmentation region. Imagine being in the rest frame of the B when a black projectile approaches at very high Lorentz γ. (Actually the ideal such projectile is a heavy ion — unmistakably black). Then it is quite reasonable to expect all the color to be driven out of the B — except for the superheavy b quark which lags behind. The projectile is a vacuum-cleaner! The color-sweep occurs, at the very least, via the Compton scatterings of each colored parton in the target by the wall of gluons.

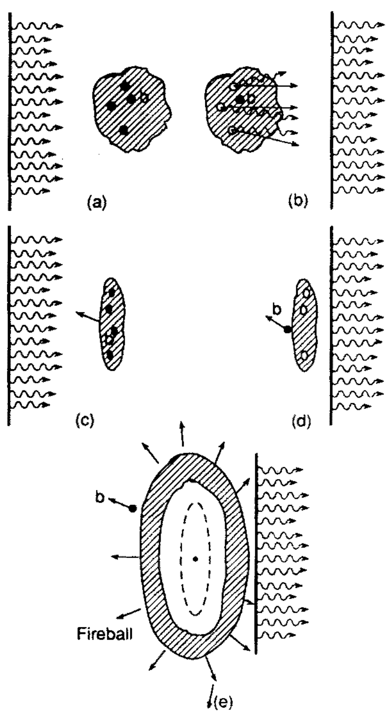

Fig. 13. (a): A cosmic-ray primary, in its rest frame, about to be hit by a high-energy air nucleus. (b): After the collision: all color is swept out of the target. (c) and (d): The same event as viewed in the rest frame of the produced fireball. (e): Somewhat later in the rest frame of the fireball — is there disoriented chiral condensate in the interior??

The products go forward at the speed of light and therefore are stuck to the wall (Fig. 13(b)).

What is left behind? In the neighborhood of the constituent quark, there is nothing at all. (In the neighborhood of the b-quark there is color separation *a la* $e^+e^- \to q\bar{q}$, but that is an inconsequential complication, not present in pp collisions). But this "nothing," which carries no memory of the valence degrees of freedom, cannot be vacuum. Therefore the "nothing" carries four-momentum P^*_μ and possesses mass M^*. If M^* is large the "nothing" must, from conservation of $(E - p_\|)$, be driven downstream when viewed in the rest frame of the struck proton. But if the "nothing" is color singlet it may become detached from the wall and drift behind it. This marvelous idea I learned from Bill Walke [16], and if realized would lead, in the cms frame, to a leading "fireball" in B direction separated from the rest of the phase-space by a rapidity-gap[3].

How does the "nothing" evolve? It should expand at the speed of light carrying most of its mass outward in a not-too-thick (??) shell of partonic matter until the energy density is appropriate for hadronization [17]. A reasonable value for the decoupling radius is when the shell is everywhere dense with one layer of pions of momentum $\langle p_T \rangle \approx 500$ MeV. Then [18]

$$M^* \sim \frac{4\pi R^2}{\pi \langle r_\pi^2 \rangle} \langle p_T \rangle \sim \langle p_T \rangle \, m_\rho^2 R^2 \,, \tag{23}$$

where $\pi \langle r_\pi^2 \rangle$ is the area of a pion, of order $4\pi/m_\rho^2 = \pi(0.6\, f)^2$. Thus

$$R \sim 0.4\, f \left(\frac{M^*}{1\ \text{GeV}}\right)^{1/2} \approx 0.7 \langle r_\pi \rangle \left(\frac{M^*}{1\ \text{GeV}}\right)^{1/2} \tag{24}$$

and for $M^* \gtrsim 50$ GeV the situation becomes rather macroscopic.

To release that much mass requires a central blackness $\nu \gg 2 \ln \sqrt{50} \sim 7.5$. The blackness in a central pp collision at the SSC has been estimated to be ~ 9.5 by Block, Cahn and Margolis [19]. Per quark it may not be as large. But one may wait for those rare configurations when all quarks shadow each other (Fig. 14). Take the probability a quark shadows the diquark to be $\sim 0.1 \langle r_q^2 \rangle / \langle r_p^2 \rangle$. This is conservative. Take the same probability for the diquark to line up, but let the diquark have 0.7 the radius of the proton. Then the probability of a fully aligned configuration is

$$0.02 \left(\frac{\langle r_q^2 \rangle}{\langle r_p^2 \rangle}\right)^2 \tag{25}$$

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[3] Actually in our example the hadronization products of the b quark could fill in the gap. But this need not happen in pp collisions.

and the fraction f of all collisions for which both protons are aligned and they hit head-on is

$$f \sim (0.02)^2 \left(\frac{\langle r_q^2 \rangle}{\langle r_p^2 \rangle} \right)^5 . \tag{26}$$

Taking $\langle r_q^2 \rangle / \langle r_p^2 \rangle = 1/9$ gives a fraction $f \sim 10^{-8}$. While this is a small number, it is still plenty large enough at hadron colliders: this is a nanobarn cross-section. The blackness ν per quark-quark collision that is needed goes down to $\gg 7.5/3 = 2.5$, an eminently reasonable value. We conclude that fireball radii in excess of $3f$ should occur *via* this mechanism. [Note that such radii are commonplace already — but it is not clear that the same initial conditions occur.]

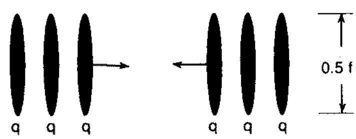

Fig. 14. The ultimate central proton-proton collision.

What is inside the fireball? At early times, the initial volume is extremely small

$$V_0 \sim \frac{4}{3}\pi(0.2\,f)^3 \sim 0.03\,f^3 \tag{27}$$

and for $M^* = 50$ GeV , the initial energy density is an impressive

$$\epsilon_0 \gtrsim 1\ \text{TeV/fm}^3 . \tag{28}$$

However things happen so fast that it may be questioned whether there exists thermal equilibrium. Presumably the gluon-wall vacuum cleaner would be effective in disturbing the fermi-sea in the neighborhood of the constituent quark. It cannot sweep it clean (that means filling the positive energy sea), but a near-degenerate Fermi gas might be produced *ab initio*. Assume each level of q and $\bar{q}$ has population one-half up to some momentum k_0 (unity seems too big; zero too small). Then we get for the number $N + \bar{N}$ of quarks plus antiquarks

$$N + \bar{N} \sim \frac{1}{(2\pi)^3} \left(\frac{4}{3}\,\pi R_0^3\right) \left(\frac{4}{3}\,\pi k_0^3\right) \frac{1}{2} \cdot 3 \cdot 3 \cdot 4 = \frac{16}{\pi}\,(k_0 R_0)^3 \tag{29}$$

and for the mass

$$M^* = \frac{3}{4}\,k_0(N + \bar{N}) . \tag{30}$$

Numerically, with $R_0 \sim 0.2\,f$, we have

$$\begin{aligned} N + \bar{N} &\sim 40 \left(\frac{k_0}{2\text{ GeV}}\right)^3 , \\ M^* &\sim 60\text{ GeV} \left(\frac{k_0}{2\text{ GeV}}\right)^4 . \end{aligned} \tag{31}$$

What happens during the expansion? With the boundary-region receding at the velocity of light, these fermion states may simply "red-shift" and ultimately annihilate, leaving behind a degenerate system with baryon number $B = 1/3(N - \bar{N})$.

Is it thinkable that this relaxation process is gentle enough to create baryonic matter, in particular strange-matter [20, 21] "nuggets"?? The probability for the initial fireball to fluctuate to, say, $N = 30$, $\bar{N} = 10$ is not too small (0.7%), and observation of even one such object ($M^* \sim 7$ GeV; $Z = 0$, $S = 7$) would be quite enough.

Slightly more thinkable is creation of disoriented chiral condensate within the fireball [22–25]. This is just vacuum with the chiral orientation tilted away from the σ direction into a π direction by an amount θ. This chiral disorientation persists until decoupling, at which time the vacuum relaxes back to the sigma direction *via* coherent emission of mainly non-relativistic pions, all of the same (Cartesian) isospin orientation. The number of anomalous pions is

$$N_{\text{coh}} = \theta^2\, f_\pi^2\, m_\pi \left(\frac{4}{3}\,\pi\, R^3\right) . \tag{32}$$

The number in a "standard" fireball is determined by the argument leading to Eq. (23)

$$N_{\text{std}} = \frac{4\pi\, R^2}{\pi\, \langle r_\pi^2 \rangle} . \tag{33}$$

Thus the ratio is

$$\frac{N_{\text{coh}}}{N_{\text{std}}} = \frac{\pi}{3}\, \theta^2\, f_\pi^2\, m_\pi\, \langle r_\pi^2 \rangle\, R \;\cong\; 0.1\,\theta^2 \left(\frac{R}{1\,f}\right) \cdot \left(\frac{\langle r_\pi^2 \rangle}{0.5\, f^2}\right) . \tag{34}$$

If the disorientation of the chiral vector is random, then the distribution of the neutral fraction $f = N_{\pi^0}/(N_{\pi^0} + N_{\pi^+} + N_{\pi^-})$ is inverse square root. There should be events where all pions are charged; others where they are all neutral. The phenomenology matches completely what is claimed [26] by the Chacaltaya–Pamir group for the Centauro/Chiron phenomena. And, in a paper speculating about possible ultracold quark-gluon plasma, Van Hove

[27] has assembled hints from accelerator data which point in the same direction.

Finally we note that the necessary condition, according to Eq. (22), for completely stopping the incident protons is $\nu \gg 2 \ln (40 \text{ TeV}) \approx 21$. This is twice what is expected for a generic central pp collision. For a collinear fluctuation such as we described, it is conceivable that ν satisfies this condition.

Thus, in addition to the leading-particle physics we described, there might be a central system which is thermalized with Landau initial conditions. The ultimate limiting case is all 40 TeV thermalized, in which case [10] the initial energy density is $\sim 10^{10}$ GeV/fm^3 and the initial temperature ~ 70 GeV. One may err by a few orders of magnitude in the fraction of energy thermalized and still expect some interesting phenomenology to occur.

It is a pleasure to dedicate this article to my friend and colleague Wiesław Czyż, who exemplifies by the way he does physics, as well as by his many important and beautiful contributions, the reason it is such a privilege to be a scientist.

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