Remarks on glueballinos (R^0 particles) and R^0 searches

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We propose a search strategy for the light R^0 (glueballino) particle suggested by Farrar in connection with the light gluino scenario. The basic idea is to moderate and stop the R^0 particles and then observe their decay to almost monochromatic π^0's, at an appropriate time delay relative to a primary collision event, where a gluino jet, likely to fragment into the R^0, was produced. This technique is optimized at colliders ($\bar{p}p$, e^+e^-), and depends on the qualitative features of the R^0 hadronic interactions which we discuss in detail.
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I. INTRODUCTION

Over the last few years a case for a light gluino and photino has been made by Farrar [1–3]. Supersymmetry breaking is expected to generate masses in the $100-10^3$ GeV range for the superpartners of the known light fermions and bosons. [The upper range is a theoretical bias and is natural if supersymmetry (SUSY) and electroweak (EW) symmetry breaking are concurrent and related—the lower limits are experimental ''bounds.''] However in certain schemes the partners of the exactly massless vector bosons (associated with unbroken $SU(3)_C \times U(1)_{EM}$ gauge symmetric), namely, the photino and gluino, obtain only small $\leqslant 0(\mathrm{GeV})$ masses via radiative (loop) correction.

Since finding a light gluino $\tilde{g}$ at existing accelerators [before the CERN Large Hadron Collider (LHC)], would have such monumental consequences, we believe that even the remote chance that this scenario is realized in nature is worth pursuing. Light gluinos modify various aspects of perturbative QCD, such as the running of $\alpha_s(Q^2)$, quarkonium, and jet physics [4,5]. It seems that these modifications cannot, at present, definitely rule out the light gluino hypothesis.

The gluino cannot, because of color confinement, directly manifest as a free particle, but rather as a constituent of a color neutral hadron the ''glueballino'' $\tilde{g}g$. This new hadron, termed $R^{(0)}$ by Farrar, is the lightest hadron with negative R parity. It is therefore strong interaction stable but decays weakly via squark exchange into a photino+hadrons:

$$R^0 \to \tilde{\gamma} + H^0 (H^0 = \pi^0, \eta^0, \pi^+\pi^-, \pi^+\pi^-\pi^0, \text{ etc.}). \quad (1)$$

(If $m_{R^0} \leqslant m_{\tilde{\gamma}}$ the R^0 would be stable altogether. This theoretically (even more) unlikely scenario is forbidden by the following consideration: after the $R^0R^0 \to$hadrons annihilation freeze out, a relic density of R^0: $n_{R^0}/n_\gamma \simeq 10^{-18}$ remains. This implies $n_{R^0}/n_p \geqslant 10^{-10}$. As we will argue below R^0 is likely to attach to heavy nuclei. The (R^0-A,Z) composites constitute new exotic ''isotopes,'' which have been excluded with very high precision [6,7].)

The R^0 decay lifetime scales as the squark mass to the fourth power and is even more sensitive to the R^0 mass (or rather $R^0-\tilde{\gamma}$ mass difference). For R^0 masses in the range

$$m_{R^0} = 1.5 \pm 0.2 \text{ GeV} \quad (2)$$

and

$$70 \text{ Gev} \leqslant m_{\tilde{q}} \leqslant 500 \text{ GeV}. \quad (3)$$

Farrar (and we) find a very wide range for R^0 lifetimes:

$$\tau_{R^0} \simeq (3\times 10^{-11} \text{ sec} - 10^{-4} \text{ sec}). \quad (4)$$

The R^0 hadron is a most striking prediction of the light gluino scheme. How come such a particle has not been discovered yet?

As correctly pointed out by Farrar [1,3], searches for SUSY particles looking for missing E_T (transverse missing energy) signals are rather insensitive to R^0's. Over most of the above lifetime range the R^0 interacts in the calorimeters and loses most of its energy there. The final decay photino would then carry a tiny ($\leqslant$ few GeV) energy and would be indistinguishable from a neutrino with similar energy. Likewise the R^0 could not have been discovered in beam dump experiments looking for penetrating particles. Finally searches for gluino jets based on their q jetlike angular distribution [5] may also be somewhat hampered by the fact that the leading hadron in these jets, namely, $\tilde{g}g \equiv R^0$, can take most (60–80 %) of the jet's energy. Unlike leading gg glueball in gluon jets which decay to multiple π's, K's hadronic, 10^{-24} sec, time scales, the R^0 are stable on jet evolution time scale. Gluino jets may therefore be less conspicuous than gluon or quark jets.

In principle neutral beam experiments with decay paths at a distance L from the production point can be used to look for R^0 decays if $L/c \simeq \gamma_{R^0}\tau_{R^0}$. Also the R^0's in the beam could be looked for via their hadronic collisions by careful measurement of time of flight and of R^0p elastic collisions kinematics which hopefully can separate the R^0 from the dominant neutron component [1,2].

We will focus on an alternative approach. The hadronic interactions of the R^0 can be utilized to moderate it. R^0 de-

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cays delayed by $\tau(R^0)$ from the ''relevant'' primary interaction (chosen on the basis of being likely to send an R^0 in the detector's direction) could then yield energetic π^0's, η^0's, and other $\pi^+\pi^-$, $\pi^+\pi^-\pi^0$ hadronic systems.

In the following we will review the basic steps and the features of the putative R^0 hadrons which control them.

II. R^0 PRODUCTION

The sensitivity of R^0 searches depends on its production rates and on the R^0 inclusive spectra, both of which we will estimate in the following.

The $\tilde{g}\tilde{g}$-gluon coupling is larger than that of the $\bar{q}q$ gluon by the ratio of the second Casimir operators for the octet and triplet representations of $SU(3)_c$:

$$g^2_{\tilde{g}\tilde{g}g}=g^2_{ggq}=[C_2(\underline{8})/C_2(\underline{3})]g^2_{\bar{q}qg}=2.25g^2_{\bar{q}qg}. \tag{5}$$

This suggests that the gluino jets, minijets, and the R^0 particles that these jets fragment into, will be more copious than charm jets and charmed particles. This is even more so when the R^0 mass is appreciably smaller than the $D-D^*$ mass: $\bar{m}_D=\frac{3}{4}m_{D^*}+\frac{1}{4}m_D=1.97$ GeV. [$m_{R^0}\leqslant\bar{m}_D$ is assumed to always be the case]. It was found that at the tevatron collider 10% of all jets had $D+D^*$ [8]. We expect at least as many R^0's containing jets.

Unlike the case of $B^{(*)}$ mesons, whose complete production pattern can be predicted from perturbative QCD, a substantial fraction of D's may be produced ''softly'' with no separate $c\bar{c}$ jets. This is *a fortiori* the case for the R^0 particle. If this soft component can be described by the phenomenological fit

$$\frac{d\sigma}{dp_T}\simeq e^{-6\sqrt{p_T^2+m^2}}, \tag{6}$$

we would have (for $m_{R^0}\simeq 1.5$ GeV) a fairly large ratio:

$$\sigma(R^0)/\sigma(D+D^*)=\frac{m_{R^0}e^{-6m_{R^0}}}{\bar{m}_D e^{-6m_D}}\simeq 16. \tag{7}$$

Unfortunately $10-10^3$ more neutrons and neutral kaons (K_0^L specifically) are produced and provide a strong background for interacting (and decaying) R^0's, respectively. The R^0 particles are produced with larger transverse momenta than the latter n's, K^0's and the other light hadrons. This feature is particularly true for particles produced in the forward ($\bar{p}$) and backward (p) fragmentation region and in the central rapidity plateau (in the $\bar{p}p$ c.m.s.). This in turn causes a broader angular distribution of the produced R^0 particles:

$$\Theta_R\simeq\frac{p_T(R)}{p(R)}>\frac{p_T(n)}{p(n)}\left[\text{or }\frac{p_T(K_L)}{p(K_L)}\right]\simeq\Theta(n)[\text{or }\Theta(K_L)]. \tag{8}$$

In particular many of the neutrons and some K_0's may be ''leading'' particles, in the fragmentation of the proton projectile in a fixed target experiment. These will than have both small p_T's and large laboratory energies and hence very small angles.

The neutral beams in various fixed target experiments, obtained by strong collimation in almost exactly the original proton beam direction, will therefore be further strongly enriched in neutrons and K_0^L. It is thus particularly difficult to look for the decays or interactions of the relatively rare putative R^0 particles in such beams.

III. THE LIFETIME OF R^0

The most important parameter for determining optimal search strategies for the R^0 particle is its lifetime τ_{R^0}. Indeed as τ_{R^0} spans the six to seven decade range indicated in Eq. (5) its very evolution as it propagates in matter will drastically change. For $\tau_{R^0}\leqslant 10^{-9}$ sec it is likely to decay within a meter or so from the collision point and manifest via a missing p_T signal. If $\tau_{R^0}\geqslant 10^{-8}$ sec then it is likely to become nonrelativistic and further lose energy by elastic R^0-nuclei collisions before decaying. Let us therefore briefly recall why these lifetimes are naturally expected (if m_{R^0} is indeed in the $m_{R^0}=1.5\pm 0.2$ GeV range prescribed).

The gluino decay $\tilde{g}^0\to\bar{q}q\tilde{\gamma}$ with $\bar{q}q=\bar{u}u$ or $\bar{d}d$, proceeds via the squark exchange diagram. If the gluino was a physical particle with mass $m_{\tilde{g}}$ we could use this Feynman diagram to compute directly the $\tilde{g}$ decay rate. By comparing with the W^- exchange in $\mu\to\nu_\mu\nu_e e$ decay we have

$$\frac{\Gamma(\tilde{g})}{\Gamma(\mu)}=\frac{\alpha_{\rm QCD}(m_{\tilde{g}})\alpha_{\rm em}}{(\alpha_{\rm weak})^2}\left(\frac{m_W}{m_{\bar{q}}}\right)^4\frac{5}{3}\left(\frac{m_{\bar{q}}}{m_\mu}\right)^5\frac{r(m_{\tilde{\gamma}}^2/m_{\tilde{g}}^2)}{r(m_e^2/m_\mu^2)}, \tag{9}$$

where the above four factors represent coupling constants and propagator ratios, $N_c(q_u^2+q_d^2)=3.5/9=5/3$ is a charge color factor summing over $u\bar{u}$, $d\bar{d}$ states of various colors, and a phase space ratio. (In particular,

$$r(z)=1-8z+8z^3-z^4-12z^2\ \ln z \tag{10}$$

is relevant in the limit $m_{\tilde{\gamma}}\gg m_u^0, m_d^0$.) The free gluino decay rate was estimated some time ago [9].

If the gluino had a ''large,'' $m_{\tilde{g}}^{(0)}\geqslant 1.2$ GeV, bare mass and, as the charm quark in $D^+=c\bar{d}$, dominated the $R^0=\tilde{g}g$ mass then, just as in the case of D^+ decay, we could use, in a spectator approximation, $m_{\tilde{g}}\simeq m_{R^0}$ to estimate $\tau(R^0)$. Equation (9) with $\alpha_{\rm QCD}(1.5\text{ GeV})\simeq 0.3$, $\alpha_{\rm em}\simeq 0.0073$, $\alpha_W=0.034$, and $m_{\tilde{\gamma}}\ll m_{R^0}\simeq m_{\tilde{g}}$ [so that $(m_{\tilde{g}}^2/m_R^2)\simeq 1$] yields

$$\tau(\tilde{g})\simeq\tau(R^0)\simeq 1.6\times 10^{-6}\text{ sec}\left(\frac{m_{\tilde{q}}}{100\text{ GeV}}\right)^4\left(\frac{0.1\text{ GeV}}{m_{R^0}(\simeq m_{\tilde{g}})}\right)^5$$
$$\simeq 2.5\times 10^{-12}\text{ sec}(m_{\tilde{q}}/100\text{ GeV})^4, \tag{11}$$

where we used in the last step $m_{\tilde{g}}\simeq m_{R^0}\simeq 1.5$ GeV. However, in the envisioned scenario $m_{\tilde{g}}^{(0)}\ll m_{R^0}$, and Eq. (11) overestimates the R^0 decay rate. Most of the R^0 mass is due to QCD dynamic effects: the chromoelectric (and magnetic) fields, i.e., surrounding gluon cloud and some $\bar{q}q$ sea. The bare gluino therefore carries a fraction xm_{R^0} of the R^0 mass with a (differential) probability $f(x)$ normalized to

$$\int_0^1 f(x)dx=1. \tag{12}$$

Following Farrar [2] we use a ''generalized spectator'' approximation and we assume that the valence gluon and other gluons and/or $\bar{q}q$ pairs ''sail along,'' as the bare, locally coupled gluino, of ''mass'' xm_{R^0}, decays. This then yields

$$\tau(R^0)^{-1}=0.4\times 10^{12}(\text{sec})^{-4}\left[\frac{100\ \text{GeV}}{m_{\tilde{g}}}\right]^4 \times\int_0^1 f(x)x^5 r[(y/x)^2]dx. \tag{13}$$

The ratio

$$y\equiv\frac{m_{\tilde{\gamma}}}{m_{R^0}} \tag{14}$$

controls further phase-space suppression due to finite photino mass via $r[(y/x)^2]=r[m_{\tilde{\gamma}}^2/(xm_{R^0})^2]$. It is convenient to parametrize the decreasing probability that the gluino carries most of the total mass by powerlike vanishing of $f(x)$ as $x\to 1$:

$$f(x)\simeq(1-x)^{\gamma}(\gamma+1). \tag{15}$$

Even for $m_{\tilde{\gamma}}=0$ the last factor in Eq. (13) above supplies sizable suppression:

$$I_{\gamma}(0)=\int f(x)x^5dx=\frac{(\gamma+1)!5!}{(\gamma+5+1)!}. \tag{16}$$

$I_{\gamma}(0)$ is smaller than ~ 0.01 for $\gamma\geqslant 3$. Clearly $I_{\gamma}(y)\equiv\int(1-x)^{\gamma}x^5 r[(y/x)^2]<I_{\gamma}(0)$ could be arbitrarily small (or zero) once $m_{\gamma}^0\to m_{R^0}$, i.e., $y\to 1$. Following Farrar we will adopt the range of lifetimes

$$\tau_{R^0}=(10^{-10}-10^{-7})(m_{\tilde{g}}/100\ \text{GeV})^4\ \text{sec} \tag{17}$$

corresponding to $2\times 10^{-5}\leqslant I\leqslant 2\times 10^{-2}$.

In the light gluino scenario the lower experimental bounds on squark masses, which are based on missing energy searches, do not apply. The squark decays (in $\simeq 10^{-26}$ sec) to a quark and gluino. The gluino, after forming the R^0, deposits most of its energy.

For $\tau(R^0)$ estimates we will therefore take squark masses in the range

$$60\ \text{GeV}\leqslant m_{\tilde{g}}\leqslant 500\ \text{GeV}. \tag{18}$$

The lower limit is based on the modification of $R\equiv\sigma_{e^+e^-\to\text{hadrons}}/\sigma_{e^+e^-\to\mu^+\mu^-}$ measured at the CERN e^+e^- collider LEP II due to $\tilde{q}\bar{\tilde{q}}$ production (see Ref. [2], and references therein). The upper bound is a theoretical bias. Substituting in Eq. (17) we find that the wide lifetime range 3×10^{-11} sec$\leqslant\tau_{R^0}\leqslant 10^{-4}$ sec is indeed theoretically expected.

The pattern of photino couplings implies that the $\bar{q}q$ state in the $\tilde{g}\to\bar{q}q+\tilde{\gamma}$ decay is

$$q_u\bar{u}u+q_d\bar{d}d=\frac{2}{3}\bar{u}u-\frac{1}{3}\bar{d}d \tag{19}$$

or, if we include the s quark,

$$q_u\bar{u}u+q_d\bar{d}d\simeq\frac{2}{3}\bar{u}u-\frac{1}{3}\bar{d}d-\frac{1}{3}\bar{s}s. \tag{20}$$

This state overlaps more strongly with the π^0 state

$$\pi^0=(\bar{u}u-\bar{d}d)/\sqrt{2}$$

than with η^0:

$$\eta^0\simeq(\bar{u}u+\bar{d}d-\bar{s}s)/\sqrt{2}.$$

In the absence of any dynamical ''form factor'' effects this would enhance the $R^0\to\pi^0\tilde{\gamma}$ over $R^0\to\eta^0\tilde{\gamma}$ by at least a factor of 9, i.e. (see Ref. [2]),

$$\Gamma(R^0\to\pi^0\tilde{\gamma})/\Gamma(R^0\to\eta^0\tilde{\gamma})\simeq 9. \tag{21}$$

However, if $m_{R^0}-m_{\tilde{\gamma}}\simeq 1.4\pm 0.3$ is on the high side of its range, it would appear that higher multiplicity final states $R^0\to\pi^+\pi^-\tilde{\gamma}$, $\pi^+\pi^-\pi^0\tilde{\gamma}$ would be preferred over $R^0\to\pi^0,\eta^0,+\tilde{\gamma}$. A crude estimate for the probability of the two body decay modes can be inferred from

$$B(D^+\to K_{\rho}^{\pi})\simeq^{3\%}_{7\%}\simeq B(R^0\to\tilde{\gamma}\pi^0). \tag{22}$$

We will assume that

$$B(R^0\to\tilde{\gamma}_{\eta^0}^{\pi^0})=0.1\pm 0.05. \tag{23}$$

It should be emphasized however that in a recent paper (the last entry in Ref. [2]) Farrar points out that approximate C invariance is likely to supress the two-body decay modes $R^0\to\tilde{\gamma}\pi^0(\eta^0)$ relative to $R^0\to\tilde{\gamma}\pi^+\pi^-$, etc.

IV. HADRONIC R^0 INTERACTIONS

An R^0 hadron propagating in matter scatters from the nuclei. These interactions are elastic if the kinetic energy of the R^0 is below the binding energy of nucleons

$$T_{R^0}\leqslant\text{BE}\leqslant 10\ \text{MeV} \tag{24}$$

and quasielastic (i.e., nucleons knock out) if

$$10\ \text{MeV}\leqslant T_{R^0}\leqslant 0.5\ \text{GeV}. \tag{25}$$

Finally at higher energies

$$T_{R^0}>0.5\ \text{GeV} \tag{26}$$

we have, in addition to the (quasi) elastic R^0N scatterings, also genuinely inelastic scattering involving production of pions, kaons, etc. Most of these interactions are soft, at least in the t channel, and cannot be reliably calculated via perturbative QCD. Yet a qualitative understanding of these interactions and, in particular, the degree to which these differ

from pion, kaon, and nucleon interactions is required, if we wish to estimate the evolution of energetic (say $E_{R^0} \simeq P_{R^0} \simeq 20$ GeV) R^0's in matter.

At very high energies it is believed that meson and nucleon scattering are dominated by the exchange of (two or more) gluons. At lower (Laboratory) energies $E_{\rm lab} \leqslant 10$ GeV quark-antiquark annihilations and exchange play a dominant role. In pre-QCD nomenclature these are the ''Pomeron'' and mesonic ''Regge trajectories'' exchanges, respectively, with the corresponding cross section

$$\sigma_{\rm pom \simeq glue\text{-}ex} \simeq \text{const} \times (\text{or } \ln W), \quad W \to \infty, \tag{27}$$

$$\sigma_{\rm Regge \simeq quark\text{-}ex} \simeq \frac{\text{const}'}{W}, \quad W \to \infty. \tag{28}$$

What are the corresponding contributions for the case of R^0N scattering? The fact that gluino (and gluon) have stronger gluonic couplings than quarks [Eq. (5)] suggests that the gluonic exchange (Pomeron) part is enhanced there. Indeed the contribution of the chromomagnetic-electric cloud energy

$$I_{R^0} = \int d^3r (E^2 + B^2)[\text{in } R^0 \text{ state}] - \int [\text{in vac}] \simeq m_{R^0} - m_{\tilde{g}}^0 \tag{29}$$

makes up most of the R^0 mass. I_{R^0} is larger than the corresponding integral in a meson or even in a nucleon. This conforms to the enhanced contributions of gluonic cloud interaction to R^0-N scattering.

On the other hand we have no analogue of the quark annihilation or exchange diagrams for R^0-N scattering. (We could have only gluon exchange analogies of these diagrams and even these only for, say, R^0R^0 scattering.)

Most of the M-N and NN scatterings for $E_{\rm lab} \leqslant 10$ GeV is not only merely accounted for by Reggeon exchange. There are further detailed hadronic models involving pion and other meson exchanges for the specific reaction mechanism in various inelastic channels. Since to lowest order there are no such pion (meson) exchanges contributing to R^0-N scattering we expect the latter to be smaller than M-N scattering, particularly for $E_{\rm lab}(R^0) \leqslant 10$ GeV. Further arguments suggest that

$$\sigma(R^0N) \leqslant \sigma(MN) \tag{30}$$

holds not only for $E_{\rm lab}(R^0) \leqslant 10$ GeV, but also at higher energies. The issue of high energy MN-NN cross sections is not completely settled. It is possible that ''soft Pomeron'' contributions due to multiperipheral pion exchange type models make up a sizable part of the total M-N cross section even at lab energies as high as 20–40 GeV.

A closely related issue is the ambiguity of which meson nucleon scattering cross section should we compare $\sigma(R^0N)$ with. These MN cross sections depend on whether we use

$$M = \pi = u\bar{d}: \sigma_{\rm tot}(\pi^+) \simeq 30 \text{ mb}, \tag{31}$$

$$M = K = u\bar{s}: \frac{1}{2}[\sigma_{\rm tot}(K^+p) + \sigma_{\rm tot}(K^-p)] \simeq 20 \text{ mb}, \tag{32}$$

or

$$M = \phi = s\bar{s}: \sigma_{\rm tot}(\phi) = 10 \text{ mb}. \tag{33}$$

[The total ϕN cross section in Eq. (33) is inferred from vector dominance in high energy ϕ photoproduction. Estimates of $\sigma_t(\rho p)$ or $\sigma_t(wp)$ done in a similar way for ρ, w photoproduction yield $\sigma(\rho^0N) \sim \sigma(w^0N) \simeq \sigma(\pi^\pm N)$ as expected for the ρ and w which, similar to the pion π, are made of light $(\bar{u}u \pm \bar{d}d)$ quarks.]

The clear cut systematics

$$\sigma_{\pi N} > \sigma_{KN} > \sigma_{\phi N} \tag{34}$$

is somewhat puzzling from the point of view of the gluon exchange model as the QCD couplings are flavor independent and $g^2_{\bar{u}ug} = g^2_{\bar{d}dg} = g^2_{\bar{s}sg}$. One may try and argue that the slightly larger constituent mass of the s quark ($\sim$500 MeV) as compared with $\sim$350 MeV for the u,d quarks, can cause, via reduced mass effects, the sizes and cross sections to differ. Even if correct, this explanation undermines our original argument for an enhanced gluon exchange contribution to the R^0N asymptotic cross section in the first place. The ''constituent gluon'' mass, if this is indeed a useful concept, and *a fortiori* that of the gluino, are presumably significantly higher (say, m_g, $m_{\tilde{g}} = 0.7$, 0.9 GeV), than those of the s quark. The naive modeling of nonrelativistic $\tilde{g}$-g states would suggest that these larger constituent masses (and stronger attractive forces) make a much smaller state. The smaller length of the ''color dipole'' would then compensate for the $\sim\sqrt{9/4} = 3/2$ larger color charges. The systematics (34) could reflect mainly the weaker contributions of pion exchange in KN and (even more) in ϕN scattering, suggesting reduced R^0N cross sections as well.

One final piece of evidence for small R^0-nucleon cross section comes from a recent HERA experiment where a small ''Pomeron-nucleon'' total cross section has been inferred. This could be relevant to our discussion if indeed the Pomeron is related to two gluon states.

In the ensuing discussion of R^0 evolution in matter we will still allow a wide range of variation for the inelastic R^0-N cross section

$$40 \text{ mb} \sim 1.4\sigma^I(NN) \geqslant \sigma^I(R^0N) \geqslant 0.7\sigma(\phi N) \simeq 7 \text{ mb}, \quad \text{for } T_{R^0} \geqslant 0.5 \text{ GeV} \tag{35}$$

and similarly for the elastic cross section

$$15 \text{ mb} \geqslant \sigma^{El}(R^0N) \geqslant 3 \text{ mb}, \quad \text{for } T_{R^0} \geqslant 0.5 \text{ GeV}. \tag{36}$$

We should keep in mind, however, that the lower values are theoretically more favorable. Another crucial parameter controlling the hadronic shower initiated by an energetic R^0 is the fraction of the initial R^0 momentum $\overline{x_F}(R^0)$ carried on average by the final R^0 in an inelastic reaction:

$R^0(\text{momentum } p \gg m_{R^0}) + N(\text{at rest})$

$$\rightarrow N' + n \text{ pions} + R^0(p_f). \tag{37}$$

In proton-proton collisions a final ''leading'' nucleon (i.e., proton or neutron) carries on average $\overline{x_F}(p) \simeq 1/2$. This presumably reflects the tendency of a ''diquark'' in the initial nucleon to reappear as a constituent diquark in the final nucleon.

A key to the evolution of R^0 in matter is that the initial gluino in R^0 $(=\tilde{g}g)$ must be conserved in all collisions. This implies that in any reaction initiated by an R^0 particle there must be an outgoing R^0 particle containing the initial $\tilde{g}$. Since the R^0 is likely to contain also the ''initial constituent gluon,'' this leading R^0 may carry even a larger fraction of the incoming momentum than an outgoing nucleon

$$\overline{x_F}(R^0) \simeq 0.65 \pm 0.15. \tag{38}$$

Indeed a large $\overline{x_F}$ is also consistent with a ''hard fragmentation'' of gluino jets into a ''strongly leading'' R^0 carrying a large fraction of the gluino jet momentum. Such a hard $\tilde{g} \rightarrow R^0$ fragmentation may be required in order to avoid the exclusion of the light gluino hypothesis by the analysis of four jet LEP events (see Ref. [5]).

At intermediate energies [Eq. (25)] the R^0 ''knocks out'' a nucleon out of the nucleus. The fraction of kinetic energy retained by R^0 can be approximately obtained just from the two body elastic R^0N collision kinematics:

$$p_1(R^0) + p_2(N) \rightarrow p_1'(R^0) + p_2'(N). \tag{39}$$

The kinetic energy of the final nucleon $T_2' = E_2' - m_N$ is given in terms of the invariant momentum transfer $t = q^2$, $q = (p_1 - p_1')$ via

$$T_2' = t/2m_N. \tag{40}$$

In terms of the $R^0 - N$ c.m.s. momentum p^* and scattering angle θ^*, t is given by

$$t = 2(p^*)^2(1 - \cos\theta^*). \tag{41}$$

Expressing p^* in terms of the initial kinetic energy of the incoming R^0, we finally find [for $(p^*/m_N)^2 < 1$] that

$$T_2' = T_1 - T_1' = \frac{2m_N m_{R^0}}{(m_N + m_{R^0})^2}(1 - \cos\theta^*)T_1. \tag{42}$$

The retained energy fraction T_1'/T_1 is thus

$$T_1'/T_1 = 1 - \frac{2m_N m_{R^0}}{(m_N + m_{R^0})^2}(1 - \cos\theta^*), \tag{43}$$

so that on average we have for $m_{R^0} = 1.5$ GeV

$$\langle T_1'/T_1 \rangle_{R^0 - N} = 1 - \frac{2m_N m_{R^0}}{(m_N + m_{R^0})^2}(1 - \overline{\cos\theta^*})$$
$$\simeq 1 - 0.48(1 - \overline{\cos\theta^*}). \tag{44}$$

If the c.m.s. angular distribution is isotropic then $\overline{\cos\theta^*} = 0$. This indeed will be the case at the lower end of the range considered where the scattering is mainly in S wave. If we allow a small admixture of P waves then

$$\overline{\cos\theta^*} = 2(\sin\delta_1/\sin\delta_0), \tag{45}$$

where δ_0, δ_1 are the S, P wave phase shifts. Assuming

$$0 \leqslant \sin\delta_1/\sin\delta_0 \leqslant 0.2 \tag{46}$$

we will have

$$\langle T_1'/T_1 \rangle = 0.62 \pm 0.1. \tag{47}$$

The ratio of the corresponding momenta or velocities is then

$$p_1'/p_1 \simeq \beta_1'/\beta_1 \simeq \sqrt{T_1'/T_1} \simeq 0.79 \pm 0.16. \tag{48}$$

Finally we address the region of slow ($T_{R^0} \leqslant 10$ MeV) R^0 particles, which should be discussed in terms of R^0-nucleus (rather than R^0-nucleon) interactions.

The evolution of R^0 particles in this stage critically depends on the existence of a bound $(R^0 - A, Z)$ state of the R^0 and various (Fe, Ca, S, Si, and O) nuclei and on the probability that such bound states form in reactions such as

$$R^0 + (A,Z) \rightarrow \begin{Bmatrix} [R^0(A-1,Z)] + n \\ [R^0(A-1,Z-1)] + p \end{Bmatrix}. \tag{49}$$

For the long wavelength, slow, R^0 the individual nucleons will effectively merge and we expect that its interaction with the nucleus can be described via a smooth, effective, potential:

$$V_{R^0}(r) = V_\infty \frac{\rho(r)}{\rho(\infty)} \equiv V_0 \frac{\rho(r)}{\rho(\infty)}, \tag{50}$$

where $\rho(r)$ is the nuclear density and $\rho(\infty)$, and $V_\infty (= V_0)$ refers to an ideal case of infinite uniform nuclear matter.

To simplify our estimates we will approximate also finite nuclei by a ''square-well'' potential,

$$V_{R^0-(A,Z)}(r) = V_0 \Theta[R(A,Z) - r], \tag{51}$$

with r the distance of R^0 from the center of the nucleus and $R(A,Z)$ the nuclear radius:

$$R(A,Z) \simeq 1.2A^{1/3} \text{ Fm}. \tag{52}$$

If the momentum of the incoming R^0,

$$p = \sqrt{2m_{R^0}T_{R^0}}, \tag{53}$$

is sufficiently low so that

$$pR(A,Z) \leqslant 1 - 2, \tag{54}$$

we may assume that the R^0 nuclear scattering is dominated by S wave. The last condition is satisfied for the nuclei considered with $56 \geqslant A \geqslant 16$ only if $T_{R^0} \leqslant 10$ MeV, which is indeed assumed here. Also Eq. (54) justifies the smearing of individual potentials of R^0 interaction with the various nucleons.

The S wave phase shift is given by

$$\delta_0 = -pR + \arctan\left\{ \frac{p}{\sqrt{p^2+2mV_0}} \tan(\sqrt{p^2+2mV_0}R) \right\}. \tag{55}$$

The single parameter $V_0 = V_{R^0}$ summarizes our ignorance of R^0-nuclear interactions. We make the reasonable assumption that

$$V_{R^0} \leqslant |V_\Lambda| \simeq 50 \text{ MeV}, \tag{56}$$

where V_Λ is the attractive potential seen by a Λ particle in a hypernucleus. Indeed the R^0, just like the Λ particle, is not affected by the Pauli principle and can migrate towards the inner shells. However, the R^0-N interactions, lacking meson (and in particular pion) exchanges, are expected to be weaker.

Indeed π exchange diagrams contribute via three body interactions, to the Λ-nuclear binding. We have no analogue of this for the R^0 case. As we will argue below (towards the end of this section), the ''R^+'' companion of the R^0 particle is likely to be heavy. In particular

$$m_{R^+} - m_{R^0} \geqslant m_\pi \tag{57}$$

as compared with

$$m_\Sigma - m_{\Lambda^0} \simeq 70 \text{ MeV} \tag{58}$$

suppressing the three body putative diagrams and V_{R^0}.

The following values of V_{R^0}, consistent with the bound $|V_{R^0}| \leqslant 50$ MeV, lead to qualitatively different behaviors of R^0-nuclear scattering in the $T_{R^0} \leqslant 10$ MeV range.

(i) V_{R^0} is positive or negative, but very small:

$$|V_{R^0}| \ll T_{R^0} \simeq 10 \text{ MeV}. \tag{59}$$

(ii) V_{R^0} is positive (repulsive) and

$$V_{R^0} \simeq T_{R^0} \simeq 10 \text{ MeV}. \tag{60}$$

(iii)* V_{R^0} is negative (attractive) and of magnitude

$$2-4 \text{ MeV} \leqslant |V_{R^0}| \leqslant 10 \text{ MeV}. \tag{61}$$

(iv)* V_{R^0} is negative and

$$50 \text{ MeV} > |V_{R^0}| > 10 \text{ MeV}. \tag{62}$$

(v) V_{R^0} is positive and

$$50 \text{ MeV} > |V_{R^0}| \leqslant 10 \text{ MeV}. \tag{63}$$

We will next discuss each of these cases.

(i) Expanding the expression (55) for $\delta_0(p)$ in the small quantities mV_0R^2 and mV_0/p^2 we find

$$\delta_0^{(\mathrm{i})} \simeq -\frac{mV_0}{p^2}\left(\tan pR - \frac{pR}{\cos^2 pR}\right). \tag{64}$$

Thus for either sign of V_0 we obtain a fairly small R^0 nuclear cross section:

$$\sigma^{\mathrm{i}}(R^0 - \text{nucleus}) \simeq \frac{4\pi(mV_0)^2}{p^6}, \quad \frac{d\sigma}{d\Omega} = \text{const.} \tag{65}$$

(ii) In this case δ_0 can be appreciable:

$$\sigma^{(\mathrm{ii})}(R^0 - N) \simeq \frac{4\pi}{p^2} \simeq \pi(R_{A,Z})^2, \quad \frac{d\sigma}{d\Omega} = \text{const.} \tag{66}$$

(iii)* The cases marked with asterisks are distinguished by two important features. First we are guaranteed to have one (or more) bound states. Indeed the condition for having at least one S wave bound state

$$\sqrt{2mV_0}R(A,Z) \geqslant \pi/2 \tag{67}$$

amounts to

$$V_{R^0} \geqslant 4-2 \text{ MeV} \tag{68}$$

for nuclei in the oxygen-iron range, i.e., for

$$16 \leqslant A \leqslant 56. \tag{69}$$

Second we will argue later that cases (iii) or (iv) are the most likely in any event. For case (iii) we expect enhanced S wave scattering cross section, i.e.,

$$\sigma^{(\mathrm{iii})} \simeq \sigma^{(\mathrm{iv})} \simeq \frac{4\pi}{p^2}. \tag{70}$$

However, the small binding energy ($|\epsilon| \leqslant |V_0|$) may kinematically disallow (or suppress) the actual R^0 capture reaction (49) since there may not be sufficient energy for knocking a nucleon out. The latter requires in particular that

$$T_{R^0} + |\epsilon| \geqslant \text{BE of nucleon} \sim 8 \text{ MeV}. \tag{71}$$

(iv)* In this case we could have several S wave (and perhaps even P wave) bound states. However the most important point is that Eq. (71) is now satisfied, sometimes with a wide ''margin.''

The R^0 capture, reaction (49), leading to the formation of the $R^0-(A-1)$ bound state, becomes then exothermic and the capture cross section is enhanced by the ratio of outgoing nucleon momentum and the incoming R^0 momentum:

$$\frac{p_N(\text{final})}{p_{R^0}} \simeq \sqrt{\frac{T_{R^0} + |\epsilon| - \text{BE}}{T_{R^0}} \frac{m_n}{m_{R^0}}} \geqslant 1. \tag{72}$$

The total scattering cross section in this case $\sigma_{\mathrm{tot}}^{(\mathrm{iv})}$ is larger than $\sigma_{\mathrm{tot}}^{(\mathrm{v})}$. The last case corresponds to a flipped (repulsive) potential: $V^{(\mathrm{v})} = -V^{(\mathrm{iv})}$. Hence (see Sec. V)

$$\sigma_{\mathrm{tot}}^{(\mathrm{iv})} \geqslant \sigma_{\mathrm{el}}^{(\mathrm{v})} \simeq \pi R_{(A,Z)}^2. \tag{73}$$

The kinematic enhancement of Eq. (72) suggests that the capture component in the total cross section

$$\sigma_{\mathrm{tot}}^{(\mathrm{iv})} = \sigma_{\mathrm{el}}^{(\mathrm{iv})} + \sigma_{\mathrm{cap}}^{(\mathrm{iv})} \tag{74}$$

is appreciable.

(v) In this case the nucleus acts as a hard sphere. Furthermore $pR(A,Z)>1$, and we may use the classical cross section

$$\sigma^{(\mathrm{v})}=\pi R^2_{(A,Z)}. \tag{75}$$

For elastic R^0-nucleus collision the energy fraction carried by the outgoing R^0 is given by Eq. (43), but with $m_N \to m(A,Z)\simeq Am_N$. Thus the average kinetic energy fraction retained

$$\langle T_1'/T_1\rangle_{R^0\ \mathrm{nucleus}}\simeq 1-\frac{2uA}{A^2+2uA+u^2}(1-\overline{\cos\ \theta^*}), \tag{76}$$

with

$$u=\frac{m_{R^0}}{m_N}=1.6\pm 0.3 \tag{77}$$

is significantly closer to one than that in the case of elastic R^0 nucleon (a feature well known from neutron moderation). Equation (76) implies a corresponding ratio for outgoing and incoming R^0 velocities:

$$\langle \beta_1'/\beta_1\rangle_{R^0\ \mathrm{nucleus}}=1-\frac{uA}{A^2+2uA+u^2}(1-\overline{\cos\ \theta^*}). \tag{78}$$

We next proceed to estimates of V_0, making first an argument for negative, i.e., attractive V_0. Approximating $R^0 \simeq \tilde{g}+$gluons, i.e., neglecting additional $\overline{qq}$ constituents, there are no Pauli exclusion repulsive effects in R^0N scattering. In fact $\overline{qq}$ annihilations and quark exchanges, i.e., ordinary mesonic contributions to R^0-N scattering, repulsive or attractive, are altogether absent. The only interactions are gluonic exchanges, i.e., the QCD analogue of ''Van der Waals type'' interactions. Color confinement implies a finite lowest mass $m^{(0)}_{\text{''}2g\text{''}}>0$ in the gluon exchange channel. This will modify the $r^{-6}-r^{-7}$ potentials derived in the perturbative two massless gluon exchange approximation by an exponential factor

$$e^{-m^0_{2g}r}. \tag{79}$$

We conjecture, however, that the generic property of the $2\gamma,2g$ perturbative exchange potentials, namely, their almost universal attractive character, will survive in the full-fledged QCD treatment of R^0N interactions.

In the perturbative approximation the ''Casimir-Polder'' interaction $V^{(0)}_{2g}$ between two hadrons A and B is proportional to the following combination of electric and magnetic polarizabilities of A and B [10]:

$$C_{AB}=-[\alpha_E^{(A)}\alpha_E^{(B)}+\alpha_M^{(A)}\alpha_M^{(B)}-(\alpha_E^{(A)}\alpha_M^{(B)}+\alpha_M^{(A)}\alpha_E^{(B)})]. \tag{80}$$

For A and B hadrons which are ground states in the respective channels, second order perturbation in an $E(B)$ background fields, say, $\Delta_E^{(\mathrm{i})}\simeq\alpha_E^{(\mathrm{i})}E^2$, imply negative α_E^{i} (or α_M^{i}). Hence C_{AB} is guaranteed to be negative once at least one of the following holds: (i) Chromoelectric effects dominate: $\alpha_E^{(\mathrm{i})}\gg\alpha_M^{(\mathrm{i})}$; (ii) $\alpha_E^{(A)}=\alpha_E^{(B)}$ and $\alpha_M^{(A)}=\alpha_M^{(B)}$, or at least

$$\alpha_E^{(A)}/\alpha_M^{(A)}=\alpha_E^{(B)}/\alpha_M^{(B)}. \tag{81}$$

Hadronic constituents are generally more relativistic than valence atomic electrons so that chromomagnetic effects are more important. Also three and more gluon exchange are more important relative to two photon exchange. Thus the attractive character of the R^0-N gluonic interactions is not obvious. We proved [11] via rigorous ''QCD inequalities'' techniques that the low energy interaction between ground state pseudoscalars

$$M_{ab}=\text{''}\overline{q_a}\gamma_5 q_b\text{''},\quad M'_{cd}=\text{''}\overline{q_c}\gamma_5 q_d\text{''} \tag{82}$$

is always attractive if

$$m_a=m_c$$

and

$$m_b=m_d \tag{83}$$

but all flavors a,b,c,d are different. The latter condition does indeed exclude any q exchange or $\overline{qq}$ annihilation so that this MM' attraction refers specifically to the gluonic interactions of interest. We believe that the same result holds for the interaction between any two hadrons of similar color structure, reminiscent indeed of condition (81) for the perturbative case.

The $\tilde{g}g=R^0$ and $N=uud$ nucleon states are clearly different. There is, however, no obvious reason why the sign of their mutual gluonic interactions be reversed relative to the gluonic interactions for NN or R^0R^0. We will therefore assume a negative V_0. Unfortunately the magnitude $|V_0|$ is not known. We have suggested above [Eq. (56)] that $|V_0^{(R)}| \leqslant |V_0^{(\Lambda)}|\simeq 40-50$ MeV. We will argue next for the following likely lower bound:

$$|V_0^{(R^0)}|\geqslant 10\ \ \mathrm{MeV}. \tag{84}$$

We express the fact that most of the $R^{(0)}$ mass originates from QCD dynamics in the following suggestive manner:

$$1\ \ \mathrm{GeV}\leqslant m_{R^0}-m^{(0)}_{\tilde{g}}\simeq\int_{\mathrm{over}\ \ R^0}d^3\vec{r}(\vec{E}_c^2+\vec{B}_c^2). \tag{85}$$

In the presence of nuclear matter the last expression will be modified. This modification can be parametrized by the chromodielectric (magnetic) constants shifting away from unity. Assuming that $\int\vec{E}^2$ dominates, we have for the rest energy shift, i.e., for V_0, the following expression:

$$V_0=\left(\frac{1}{\epsilon}-1\right)\int\vec{E}^2d^2\vec{r}=\left(\frac{1}{\epsilon}-1\right)(m_{R^0}-m^0_{\tilde{g}})$$
$$\simeq-(\epsilon-1)(m_{R^0}-m^0_{\tilde{g}})\leqslant-(\epsilon-1)\ \ \mathrm{GeV}, \tag{86}$$

where we used Eq. (85) and in the last step expanded in

$$\epsilon-1\ll 1. \tag{87}$$

We note in passing that $V_0<0$ amounts to the natural assumption that

$$\epsilon-1\geqslant 0. \tag{88}$$

To estimate $\epsilon-1$ we use the analogue of the familiar QED expression

$$(\epsilon-1)=\frac{4\pi}{3}\,n\,\alpha_E^{(N)} \tag{89}$$

with $\alpha_E^{(N)}$ the (chromo) electric polarizability of the nucleon and

$$n=(1.2\ \ \mathrm{Fm})^{-3} \tag{90}$$

the nucleon number density. Ordinary electric polarizability can be written as

$$\alpha_E\simeq r_0^3, \tag{91}$$

where in the last expression r_0 is an effective size [Eq. (91) also represents the classic polarizability of a conducting sphere of radius r_0]. Collecting all the above results we find

$$|V_0|\geqslant 2\left(\frac{r_0^{N,\mathrm{eff}}}{\mathrm{Fermi}}\right)^3\ \ \mathrm{GeV}. \tag{92}$$

Thus $|V_0|\geqslant 10$ MeV once the mild requirement

$$r_0^{N,\mathrm{eff}}\geqslant 0.17\ \ \mathrm{Fermi} \tag{93}$$

is satisfied.

Before concluding this section we would like to reiterate the feature which is the most important distinction between the R^0-nucleon scattering and the scattering of any other stable hadron on a nucleon.

All the hadrons with lifetimes $\geqslant 10^{-10}$ sec (with the exception of Λ^0) come in isospin multiplets:

$$\begin{pmatrix}\pi^+\\ \pi^0\\ \pi^-\end{pmatrix}\quad\begin{pmatrix}K^+\\ K^0\end{pmatrix}\quad\begin{pmatrix}\overline{K^0}\\ K^-\end{pmatrix}\quad\begin{pmatrix}p\\ n\end{pmatrix}\quad\begin{pmatrix}\Sigma^+\\ \Sigma^0\\ \Sigma^-\end{pmatrix}\quad\begin{pmatrix}\Xi^+\\ \Xi^0\\ \Xi^-\end{pmatrix}.$$

Also the Λ^0 is associated in SU(3) flavor, with the Σ multiplet from which it is split by only 70 MeV [see Eq. (58)].

In principle we could have in addition to $R^{(0)}=\tilde{g}g$ other more complex hadrons containing a gluino which could be charged as

$$R^+=\tilde{g}(u\overline{d})_8, \tag{94}$$

where as indicated the $\overline{ud}$ are in a color octet combination. However, the R^+ is definitely not an I-spin partner of R^0, but rather the lowest member of a completely different family of particles. These particles are the SUSY partners of the ''Meikton'' or ''hemaphrodite'' hadronic states, which have been speculated by various authors, namely, $g(u\overline{d})_8$, just as $R^0=\tilde{g}g$ is the SUSY partner of the glueball state gg.

To date we have no clear-cut evidence for either glue-ball or ''Meikton-hemaphrodite'' states. Theoretical prejudice and estimates within bag models [14] suggest that the glueballs are significantly lighter:

$$m^0(gg)<m^0[g(u\overline{d})_8], \tag{95}$$

which in turn suggests that also

$$m_{R^0}\equiv m^0(\tilde{g}g)<m^0[\tilde{g}(u\overline{d})_8]=m_{R^+}. \tag{96}$$

Fortunately we have also direct ''experimental'' evidence that the R^+ should be heavier than the putative light R^0 by at least m_π [see Eq. (57)]. Indeed if the reverse, namely,

$$m_{R^+}\leqslant m_{R^0}+\pi^+, \tag{97}$$

is true, then we would have a new, (almost) stable, charged hadron of mass $m_{R^\pm}\leqslant 1.6\pm 0.2$ GeV.

Its production cross section, particularly at large p_T, should considerably exceed these of the deutron d or antideutron $\overline{d}$. Such particles should have been detected for lifetimes in the range considered:

$$3\times 10^{-9}\leqslant\tau_{R^+}\simeq\tau_{R^0}\leqslant 10^{-4}. \tag{98}$$

Thus the R^0 is really conserved in hadronic collisions. In any reaction initiated by R^0 the final state will also include an R^0. In particular, any R^+ produced would decay via

$$R^+\rightarrow R^0\pi^+ \tag{99}$$

on hadronic time scales.

One final comment concerns the small nuclear effects on R^0 nuclear scattering at high and intermediate energies. These ''shadowing'' effects tend to reduce the cross section of any projectile x on a nucleus from the naive, weak coupling value

$$\sigma^{(0)}(x,A)=A\sigma(x,p) \tag{100}$$

and correspondingly prolong the interaction mean free path from

$$l^{(0)}(x,A)=\frac{1}{n\sigma(x,p)}=\frac{1}{\rho\mathrm{Av}\sigma(x,p)} \tag{101}$$

with Av=Avogadro number$=6\times 10^{23}$ and ρ the material density so that $n=\rho\mathrm{Av}$ is a nucleons number density. The shadowing effects are appreciable when

$$\sigma^0(x,A)\equiv A\sigma(x,p)\simeq\pi R^2(A,Z)=\pi A^{2/3}(1.2)^2(\mathrm{Fermi})^2 \tag{102}$$

or

$$\sigma(x,p)\geqslant A^{-1/3}\times 45\ \ \mathrm{mb}\simeq(18\text{–}12)\ \ \mathrm{mb} \tag{103}$$

for $A=16$–56.

Thus we will use for iron longer interaction path than those implied by Eqs. (86),(84):

$$l_{mfp(R^0\ \mathrm{iron})}=(10\text{–}30)\ \ \mathrm{cm} \tag{104}$$

with the lower (higher) value corresponds to $\sigma^{I}_{(R^0-N)} = 40(7)$ mb with substantial $\sim 50\%$ (no) shadow corrections.

V. THE EVOLUTION OF ENERGETIC R^0 PARTICLES IN MATTER

Energetic, $E_{\text{initial}}(R^0) \geqslant 30-40$ GeV, R^0 particles evolve as they propagate in matter in three main stages corresponding, in reverse order, to the energy ranges of Eqs. (24), (25), and (26) above. We will next try to follow the evolution in these stages, utilizing the estimates of the previous section.

Stage (a). Fast R^0's are conserved through their hadronic cascade evolution, if we neglect the presumably small probability to generate further R^0R^0 pairs in R^0N collisions with c.m.s. energies $W_{R^0N} \simeq (2E_{R^0}m_N)^{1/2} \leqslant 10$ GeV. This drastically simplifies their evolution as compared with that of other long lived neutral particles such as K^0's, which are created in secondary hadronic collisions, and neutrons, which, even neglecting secondary $\bar{N}N$ pair creation, can be knocked out of nuclei. All these processes notwithstanding the hadronic showers generated by $\bar{q}q$ jets (and also $\bar{c}c, \bar{b}b$ jets) essentially quench at $l_{\text{que}} \simeq 10$ hadronic interaction lengths inside the calorimeter or shielding materials.

We would like to argue that as they evolve through the energy range 40 GeV$\geqslant T_{R^0} \geqslant 0.5$ GeV, the R^0's will typically get further away from the initial interaction point than l_{que}, i.e., further than typical hadronic showers. Several effects contribute to this.

(i) Charged particles continuously lose energy via ionization and stop at well defined ranges, even if they have no hadronic interactions. These losses increase at low (sub GeV) energies: thus the ranges for 2 GeV p, K^+ and π^+ in iron are

$$R(p) = 16 \text{ cm}, \quad R(K^+) = 30 \text{ cm}, \quad R(\pi^+) = 60 \text{ cm}. \tag{105}$$

These ionization losses can indirectly contribute also to the slowing down of K^0's and neutrons. At $T_H \simeq 1/2-2$ GeV hadron's kinetic energies, we have appreciable cross sections for charge exchange reactions

$$K^0(\bar{K}^0) + N \to K^+(K^-) + N', \tag{106}$$

$$n + N \to p + N', \tag{107}$$

with N referring to a nucleon in the nucleus. Also in inelastic collisions at $T_H = 2-40$ GeV the incoming K^0 or neutron often ''fragments'' into a leading $K^{\pm}$ or proton. Thus roughly during half of the evolution of the hadronic showers the nucleonic or ''kaonic'' component manifests as the charged component of the isospin multiplets which slow down via ionization. As emphasized towards the end of the previous section there is no analog $R^0 \to$stable $R^{\pm}$ conversion and there are no, even indirect, ionization losses during the R^0 evolution [13].

(ii) The R^0 particles are likely to maintain higher fraction $x_F(R^0)$ of their initial momenta than nucleons or kaons.

(iii) The R^0-nuclear inelastic cross sections in general and particularly in this 40 GeV$\geqslant T_{R^0} \geqslant 1/2$ GeV range are likely to be smaller than N-N and even KN cross sections.

To simplify the treatment of R^0 evolution in stage (a) we will neglect the slow buildup of transverse momenta, namely, momenta transverse to the initial entry direction of the R^0 particle into the material. (After n collisions we have

$$\overline{|p_T^{(n)}|^2} \simeq n\Delta_0^2 \tag{108}$$

with $\Delta_0 \simeq 1/2$ GeV typical soft momentum transfer.) This leads then to a one-dimensional evolution in momentum space (or more conveniently in rapidity space $y_n = -\ln x_n$ with x_n the fraction of the initial momentum carried by the R^0 after n collisions) and in coordinate space. The coordinate is the distance r (measured along the momentum of initial R^0) from the entry point to the material, to the R^0 location. For energy-independent cross sections (particularly appropriate for R^0N scattering which are devoid of the varying Reggeon contribution) these y and r evolution decouple. In a uniform medium, assumed for simplicity, the mean interaction length for $R^{(0)}$ is then a constant, l_0. Let $g^{(n)}(r)$ be the r distribution after n collisions. By definition

$$g^{(1)}(r) \simeq e^{-r/l_0}. \tag{109}$$

We can readily derive the recursive relation

$$g^{(n+1)}(r) \simeq \int_0^r dr' g^{(n)}(r') e^{-(r-r')/l_0} \tag{110}$$

and its solution

$$g^{(n)}(r) \simeq e^{-r/l_0} r^{n-1}. \tag{111}$$

The average distance after n steps is then

$$\overline{r^{(n)}} = n l_0 = n\overline{r^{(1)}} \tag{112}$$

as expected. The rapidity space evolution will be completely analogous if we assume that after a single collision we have rapidity distribution $f^{(1)}(y) = e^{-\beta y}$ (corresponding to x^{β} in x space) so that $\overline{y^{(1)}} = 1/\beta$.

Let us denote by $\lambda^{(a)}$ the energy-momentum average reduction factor in stage (a):

$$\lambda^{(a)} \equiv \overline{x_F^{(R^0)}}. \tag{113}$$

According to Eq. (38) we take $\lambda^{(a)} = 0.65 \pm 0.15$. Starting with an initial energy $E_0 = T_0 (= 40 \pm 10$ GeV), we need then altogether $N^{(a)}$ collisions in stage (a) so as to degrade the kinetic energy to the required $T_f (= 1/2$ GeV), where

$$N^{(a)} = \frac{\ln(E_0/T_f)}{\ln(1/\lambda)} \; [\simeq \ln(80 \pm 20)/\ln(1.6 \pm 0.3) \simeq 7-18]. \tag{114}$$

(We will consistently use ( ) brackets following equations for the actual numerical estimates used in applying the equations.) The total penetration distance traveled by the $R^{(0)}$ particle in stage (a) is then on average

$$L^{(a)} \simeq N^{(a)} l_{\rm int}^{(a)} = [(12\pm 5) l_{\rm int}^{(a)} = (3\pm 2.3) \text{ iron meter equivalent}], \tag{115}$$

where we used Eqs. (114) and (104). The very large spread in Eq. (115) represents the accumulated uncertainties in $\sigma^I(R^0N)$ or $l_{\rm int}^{(a)}$ and in the elasticities $\lambda^{(a)}$ and finally in the initial R^0 energy $E^{(0)}=40\pm 10$ GeV. We believe that the values $L_{\rm min}^{(a)}=0.7$ m or $L_{\rm max}^{(a)}=5.3$ m are very unlikely. In particular to achieve $L_{\rm min}$ we need to assume $\sigma_{R^0N}^{\rm inelastic}\simeq 40$ mb [and $\overline{x_F}(R^0)\simeq 1/2$] which are too high (low), respectively, and for $L_{\rm max}$ we need $\sigma^I\simeq 7$ mb [and $\overline{x_F}(R^0)\simeq 0.8$] which appear to be too low (high), respectively. We will adopt

$$L^{(a)}=3\pm 1.5 \text{ m}. \tag{116}$$

The main conclusion which we would like to emphasize is that the putative R^0 particle travels beyond the full extent of (e.m.+hadronic) showers due to normal (non R^0) hadronic jets. Indeed the extent of the normal shower is affected also by ionization losses of the leading particles to which the R^0's are completely immune. Throughout most of the stage (a) the R^0's move with $\beta_{R^0}\sim 1$, hence its total (lab-time) duration is

$$t^{(a)}\simeq L^{(a)}/c \simeq 10^{-8}(1^{+3}_{-0.4}) \text{ sec}. \tag{117}$$

Because of time dilation effects the required proper time (and hence the lower bound on τ_{R^0} required if indeed the R^0 particle is to complete this stage) is lower than the above $t^{(a)}$. Specifically we find that

$$\tau^{(a)}=\sum_{i=0}^{N^{(a)}-1}\tau_i=\frac{l_0}{c}\sum_{0}^{N^{(a)}-1}(\lambda^{(a)})^i=\frac{l_0}{c}\left[\frac{1-(\lambda^a)^{N^{(a)}}}{1-\lambda^a}\right] \simeq (2-5)\frac{l_0}{c}\simeq (0.6-5)\times 10^{-9} \text{ sec}, \tag{118}$$

where in the last equation, τ_i, the proper time lapse between the N_a-i and N_a-i+1 collision, is shortened relative to l_0/c by $1/\gamma^{(i)}\simeq(\lambda_a)^i$ and we neglected $(\lambda_a)^{N_a}=T_f/E_0\simeq\frac{1}{80}$. Thus the modest requirement

$$\tau_{R^0}\geqslant 5\times 10^{-9} \text{ sec} \tag{119}$$

suffices to ensure that the initial R^0 particles survive stage (a).

Stage (b). Intermediate energies. Using Eq. (47) for the average fraction of kinetic energy $\lambda^{(b)}(=0.62\pm 0.1)$ retained in quasielastic R^0-nucleus collisions in stage (b) we need on average

$$N^{(b)}=\frac{\ln[T_{\rm initial}^{(b)}/T_{\rm final}^{(b)}]}{\ln(1/\lambda^{(b)})}\left(=\frac{\ln 50}{\ln 1.6}\simeq 8\right) \tag{120}$$

collisions to complete this stage. We used in (120) $T_{\rm initial}^{(b)}=0.5$ GeV, $T_{\rm final}^{(b)}=0.01$ GeV for the kinetic energies characterizing the entry to and exit from stage (b). The kinematics of the elastic R^0-nucleon collision implies that for $(d\sigma/d\Omega^*)(R^0)\simeq$const (i.e., constant differential cross section in the c.m.s. frame) we have in the lab frame for $u=m_{R^0}/m_N$ [12]

$$f_u(\cos\theta)\equiv\left.\frac{d\sigma}{d\Omega}\right|_{R^0}=\frac{1+u^2\cos 2\theta}{(1-u^2\sin^2\theta)^{1/2}} \left(\simeq\frac{1+(2.5\pm 0.6)\cos 2\theta}{[1-(2.5\pm 0.6)\sin^2\theta]^{1/2}}\right). \tag{121}$$

Let the initial R^0 direction be $\hat{n}_0$, and $\hat{n}_k$ be the direction after k successive scatterings. In $f_u(\cos\theta)$ the $\cos\theta$ argument refers to $\hat{n}_k\cdot\hat{n}_{k-1}$. The distribution after k steps will be

$$f^{(k)}(\cos\theta)\equiv f^{(k)}(\hat{n}_k\cdot\hat{n}_0)=\int d\hat{n}_1 d\hat{n}_2\cdots d\hat{n}_k f_u(\hat{n}_0\cdot\hat{n}_1)\times f_u(\hat{n}_1\cdot\hat{n}_2)\cdots f_u(\hat{n}_{k-1}\cdot\hat{n}_k). \tag{122}$$

Using the completeness relation

$$\frac{1}{4\pi}\int d\hat{n}' P_l(\hat{n}\cdot\hat{n}')P_{l'}(\hat{n}'\cdot\hat{m})=\frac{1}{2l+1}\delta(l,l')P_l(\hat{n}\cdot\hat{m}) \tag{123}$$

we find that if

$$f_u(\cos\theta)=\sum_{l=0}^{\infty} f_l(u)(2l+1)P_l(\cos\theta) \tag{124}$$

then

$$f^{(k)}(\cos\theta)=\sum_{l=0}^{\infty}[f_l(u)]^k(2l+1)P_l(\cos\theta). \tag{125}$$

We will be particularly interested in the net extra displacement of R^0, namely, the extra displacement along the initial R^0 direction ($\hat{n}_0$) accumulated during stage (b): $\Delta L^{(b)}=l_0\Sigma_{k=1}^{N^{(b)}}\cos(\widehat{n_k n_0})$. On average it will be

$$\overline{\Delta L^{(b)}}=l_0\sum_{k=1}^{N^{(b)}}\overline{(\hat{n}_k\cdot\hat{n}_0)}=l_0\sum_{k=1}^{N^{(b)}}(f_1)^k=\frac{l_0 f_1}{1-f_1}(1-f_1^N) \tag{126}$$

with

$$f_1(u)=\int d(\cos\theta)f_u(\cos\theta)P_1(\cos\theta) =C\int_a^1 dx\, x\,\frac{1+u^2(2x^2-1)}{[(1-u^2)+u^2x^2]^{1/2}}, \tag{127}$$

where

$$a\equiv\frac{(-1+u^2)^{1/2}}{u}$$

and C is a normalization constant:

$$C^{-1}=\int_a^1 dx\,\frac{1+u^2(2x^2-1)}{[(1-u^2)+u^2x^2]^{1/2}}. \tag{128}$$

The transverse displacement is given by a ''random walk''

$$(\Delta_T^{(b)})^2=(l_0^{(b)})^2\sum^{N^{(b)}} \sin^2(\widehat{n_k\ n_0})\left(\leqslant\frac{2}{3}N^{(b)}l_0^2\right) \quad (129)$$

using the above with $l_0^{(a)}\simeq l_0^{(b)}\simeq 10-30$ cm, $u=1.6\pm 0.3$, $N^{(b)}=8^{+6}_{-3}$ we estimated

$$\Delta L^{(b)}\simeq 0.4-1.3 \text{ m}, \quad (130)$$

$$\Delta_T^{(b)}\simeq 0.5\pm 0.25 \text{ m}. \quad (131)$$

Finally the R^0 motion during most of stage (b) is nonrelativistic. The velocity after the nth collision is therefore

$$\beta_n=\beta_0(\sqrt{\lambda^{(b)}})^n. \quad (132)$$

The total time ($\simeq$proper time) consumed in this stage is

$$\Delta t^{(b)}=\Delta\tau^{(b)}=\sum \Delta t_n^{(b)}=\frac{l_0}{\beta_0 c}\frac{1-(\lambda_a)^{[N^{(b)}/2+1/2]}}{1-(\lambda^{(b)})^{1/2}}$$
$$(\simeq 1.5-5\times 10^{-8} \text{ sec}) \quad (133)$$

with $\beta_0\simeq 0.8$ the initial velocity corresponding to $T_{R^0}=0.5$ GeV.

Thus if

$$\tau_{R^0}\geqslant\tau^{(a)}+\Delta\tau^{(b)}\simeq(2.5^{+6}_{-0.5})\times 10^{-8} \text{ sec} \quad (134)$$

the R^0 particle would, on average, live through the end of stage (b). Its distance from the entry point is then

$$L_{ab}=L^{(a)}+\Delta L^{(b)}\simeq 4\pm 1.6 \text{ m} \quad (135)$$

with a transverse spread of

$$\Delta_T=0.5\pm 0.25 \text{ m}. \quad (136)$$

Stage (c): Our discussion in Sec. IV suggests that once $T_{R^0}\leqslant 8$ MeV (or the R^0 velocity satisfies $\beta\leqslant\beta_0=0.1$) the R^0 particles will be quickly captured into $R^0-(A,Z)$ bound states. This would be effectively ''fix'' the R^0 and prevent further diffusion away from the R^0 location it reaches by the end of stage (b) at the (L,Δ_T) values quoted above [Eqs. (135), (136)].

Even if the R^0 particle stays unbound for the duration of its lifetime τ_{R^0}, the extra diffusion distance $\sqrt{N^{(c)}}l_0$ grows very slowly with τ_{R^0} as $\sqrt{\ln\tau_{R^0}}$. From Eq. (78) above we find that the time intervals between successive R^0-nuclear collisions are

$$\Delta t^{(n)}=\Delta t^0\left(1+\frac{u}{A}\right)^n\equiv\frac{l_0}{\beta_0 c}\left(1+\frac{u}{A}\right)^n. \quad (137)$$

Thus all $N^{(c)}$ collisions of stage (c) will be completed after a time

$$t^{(c)}\simeq\tau_{R^0}=\sum^{N^{(c)}}\Delta t^{(n)}\left\{\left(1+\frac{u}{A}\right)^{N^{(c)}+1}-1\right\}\frac{A}{u}$$
$$=\Delta t^{(0)}\frac{A}{u}e^{u/A\,N_c}, \quad (138)$$

where $N^{(c)}\gg 1$ is implicit. Specifically we find for $\tau_{R^0}=10^{-4}$ sec

$$N^{(c)}=\frac{A}{u}\ln\left(\frac{\tau_{R^0}}{\Delta t^0}\frac{u}{A}\right)$$
$$(\simeq 80 \text{ for } A=56, \quad l_0=15 \text{ cm}, \quad u=m_{R^0}/m_N\simeq 1.6), \quad (139)$$

where l_0 was estimated via the geometric R^0-(A,Z) cross section $\sigma=\pi R^2(A)=3.6\times 10^{-25}$ cm^2 to be

$$l_0\simeq 15 \text{ cm in iron}. \quad (140)$$

For $\tau_{R^0}=10^{-6}$ sec we find instead of Eq. (139)

$$N^{(c)}=40 \text{ (for } A=56, \quad l_0=15 \text{ cm}, \quad \Delta t_0=5\times 10^{-9} \text{ sec}). \quad (141)$$

The maximum total extra diffusion distance added in stage (c) is therefore

$$|\Delta\vec{R}^{(c)}|=\sqrt{N^{(c)}}l_0=\begin{Bmatrix}1.4 \text{ m}, & \tau_{R^0}=10^{-4} \text{ sec}\\ 0.9 \text{ m}, & \tau_{R^0}=10^{-6} \text{ sec}\end{Bmatrix}. \quad (142)$$

In passing we note that once R^0 (or a neutron) achieves a final thermal velocity ($\simeq 2$ km/sec, or $\beta_f=0.6\times 10^{-5}$) one reverts again to ordinary diffusion with $\Delta R\simeq\sqrt{\tau}$. However, even after 100 elastic collisions with iron nuclei a R^0 particle with starting velocity $\beta_0=0.1$ will on average go down only to $\beta_f=0.005$. A key observation is that if the R^0 decays at any time during stage (c) we expect its final velocity to be

$$\beta_f\leqslant 0.1=\beta_0. \quad (143)$$

This holds also if the R^0 is bound to the nucleus with binding energy $|\epsilon|\simeq 8$ MeV. The Fermi momentum of the bound R^0 ($\simeq\sqrt{2|\epsilon|m_{R^0}}\simeq 0.16$ GeV) corresponds to

$$\beta_f\simeq 0.1. \quad (144)$$

Consequently the lab energy of the π^0 from the decay $R^0\to\pi^0\tilde{\gamma}$ will be rather monochromatic. Specifically

$$E_{\pi^0}=(1+\beta_f\cos\theta)E_{\pi^0}^{(0)}\equiv(1+\beta_f\cos\theta)\frac{m_{R^0}^2-m_{\tilde{\gamma}}^2+m_{\pi^0}^2}{2m_{R^0}} \quad (145)$$

satisfies (on average) $|E_{\pi^0}(\beta,\cos\theta)-E_{\pi^0}^{(0)}|/E_{\pi^0}^{(0)}\leqslant\beta_f/2=0.05$.

We conclude this section with a short discussion of the likely location of this R^0 decay event, $\vec{r}$ (decay), measured relative to the intersection point. We will consider first the longitudinal displacement r_L along the initial R^0 direction $\hat{n}_0$. It is accumulated only in stages (a) and (b) and assuming $\tau_{R^0}\geqslant 10^{-8}$ sec so that stage (b) was completed we have from Eq. (135)

$$r_L\simeq L_{(a)+(b)}\equiv L_{ab}\simeq 4\pm 1.6 \text{ m}. \quad (146)$$

The total displacement has also random diffusive ''components'' from stages (b) and (c). Thus finally the total displacement is

$$\langle \vec{r}^2 \rangle^{1/2} \simeq r \simeq \sqrt{(L_{ab})^2 + (\Delta_B^T)^2 + (\vec{R}_c)^2} (\simeq 4.5 \pm 2 \text{ m}). \tag{147}$$

And its average direction relative to $\hat{n}_0$ is

$$\hat{n}_0 \cdot \hat{r} \simeq r_L / r \simeq 0.8. \tag{148}$$

The above estimates referred to iron, and also neglected the freely traveled distance to the calorimeters and hence r is underestimated. We note that lighter (concrete rock+earth materials) may actually expedite the slowing down of R^0 particles in stage (c) because the fraction of energy loss is $\simeq 1/A$.

VI. THE PROPOSED R^0 SEARCH STRATEGY

The search strategy proposed here addresses a wide range of putative R^0 parameters. Specifically we will assume $\tau_{R^0} \geqslant 3\times 10^{-8}$ sec, $\sigma_{R^0N} < \sigma_{NN}$ for both elastic and inelastic cross sections and $\overline{x_F}(R^0) \geqslant \overline{x_F}(p) = 0.5$ for the momentum fraction retained by the leading R^0 particle in inelastic collisions. The search is optimized if $\sigma_{R^0N} \simeq \frac{1}{2}\sigma_{NN}$, $\overline{x_F}(R^0) \simeq 0.7$ and particularly when $B(\pi^0)$, the branching ratio for the $R^0 \to \pi^0 \tilde{\gamma}$ decay mode is large, say $B(\pi^0) \geqslant 20\%$.

Let us assume that a detector D of volume V_D has been installed underground at, say, the FNAL collider (or at e^+e^- colliders) at a distance $R(=5\pm 2 \text{ m})$ from the intersection point and with $\vec{R}$ perpendicular to the beam direction. This detector searches $R^0 \to \pi^0 \tilde{\gamma}$, $\eta^0 \tilde{\gamma}$ and possibly also $R^0 \to \pi^+\pi^-\tilde{\gamma}$, $R^0 \to \pi^+\pi^-\pi^0\tilde{\gamma}$ decays. The signature distinguishing such decays from background events could be an almost monochromatic π^0 with $E_{\pi^0} \geqslant 0.6$ GeV. Alternatively final states including several pions [and possibly some nucleons knocked out when the decay $R^0 \to \pi^+\pi^-\pi^0\tilde{\gamma}$ occurrs inside the nucleus, once the $(R^0 - A, Z)$ bound state formed] should have a relatively isotropic distribution with limited missing momenta $p_{\tilde{\gamma}} \leqslant m_{R^0}/4 \simeq 0.4$ GeV. Our arguments for a longer migration length for R^0's as compared with the extent of hadronic showers suggest a reduction of the neutron background at the location of the detector if it as placed at a distance r_0 given by Eq. (147) above.

The collision of such neutrons with nuclei in our detector can produce π^0, $\pi^+\pi^-$, etc., final states and constitute a background to the R^0 decays. However, only neutrons with kinetic energies $T_{R^0} >$ GeV ($T_{R^0} > 2$ GeV) are likely to produce π^0 ($\pi^+\pi^-\pi^0$) in n,N collisions. These background events will have achromatic π^0 or apparent missing momenta, often pointing in the direction of $\hat{r}$, i.e., towards the decay vertex.

One can further reduce the energetic neutrons background in our detector by excluding time ''windows'' Δt, say of magnitude $\Delta t \simeq 3\times 10^{-8}$ sec following each collision event or, if the rep cycle of the collider is short (as in the upgraded Fermilab Tevatron), the first $\Delta t \simeq 3\times 10^{-8}$ sec following relevant events with jets. (Only 1–2 GeV neutrons from such events emitted in the direction of the jets are likely to make it to the required transverse distance from the colliding beam. In typical minimal bias events the neutron along with most other particles produced in the beam's jets are forward backward peaked.) The total distance traveled by relativistic hadrons in this Δt is 9 m, and most of the energetic neutrons are likely to have interacted and slowed down by then. Yet the slowly diffusing R^0 may well have not decayed if $\tau_{R^0} \geqslant 3\times 10^{-8}$ sec. This negative time correlation will also reduce background from energetic muons emerging from the primary collision which interact in the vicinity of our detector and produce secondary neutrons. Such muons, cosmic ray muons, and charged particles background in general, can be further reduced by an anticoincidence network surrounding our detector at a distance of $\simeq 1$ m from its boundaries. Also the very strong, slow, or even thermal, ambient neutron background, causing many low energy depositions, via nuclear captures, can be reduced by using a proper neutron absorber shielding. Finally we note that having over most of the distance $\vec{r}$ between the intersection point and the detector iron magnetized in a direction perpendicular to $\vec{r}$ could also drastically reduce the neutron background by having the intermediate proton curve in the magnetic field [13].

All this notwithstanding we cannot hope to conclusively ''pin down'' R^0 particles by looking at their decays only. The basic reason is that, even if the detector is placed at an optimal distance r [see Eq. (147)] (which is *a priori* known only within $\pm 60\%$), and the R^0 particles are isotropically distributed, only a small fraction

$$f = \frac{V_D B(R^0 \to \pi^0)}{(4\pi/3) r^3} \simeq [V_D/m^3](3\times 10^{-5} - 2\times 10^{-3}) \tag{149}$$

of the R^0's will decay via the $\pi^0\tilde{\gamma}$ mode, inside the detector. (In the above estimate we used $r = 5\pm 3$ m, $B = 0.1 \pm 0.05$).

To pick out the genuine R^0 decays we correlate events in the detector with primary $\bar{p}p$ collision both temporally and directionally. Specifically we will focus on primary $\bar{p}p$ collisions with $\tilde{g}$ jets that could give rise to a sufficiently high energy R^0 that in turn could make it to our detector.

Thus we will consider those ''special'' $\bar{p}p$ events with one (or two) jet(s) of transverse energy $E_T \geqslant E_T^{(0)}$ ($\simeq 10$ GeV), and with no energetic lepton(s) (to avoid b,c jets) in the jet(s). If a light gluino exists these would be gluino jet(s) with an appreciable $\simeq 20\%$ probability. Beam crossing and one $\bar{p}p$ collision on average happens every

$$\delta_{t_H} = 3\times 10^{-6} (10^{-7}) \text{ sec}$$

in the old (upgraded) versions. However, only one in $N_H \simeq 10^6$ collisions is ''special'' in the above sense. (N_H is the ratio of total $\bar{p}p$ cross section to a perturbative estimate for the cross section of the gluino-jet production).

Hence the ''special'' collisions will be spaced, on average, by $\Delta t_H = N_H \delta t_H \simeq 3-(0.1)$ sec. If the $\Delta E \geqslant 0.6$ GeV energy deposition in our detector is due to a decay of an R^0 from a gluino jet in a particular ''special'' event then we expect that: (i) the decay should happen on average at time $\simeq \tau_{R^0}$ after the primary special collision; (ii) the jet (or at least one jet) in the primary event should point in the direction of our detector within a $\Delta\theta^0$ uncertainty with $\cos(\Delta\theta^0)$ given by Eq. (148).

Jointly these two requirements will reduce the background by

$$(\tau_{R^0}/\Delta t_H)(\Delta\Omega/4\pi)\left(\simeq\frac{3\times10^{-8}-10^{-4}}{3-0.1}\right)\left(\frac{1}{20}-\frac{1}{10}\right)$$
$$\simeq(5\times10^{-10}-10^{-4}), \qquad (150)$$

where

$$\Delta\Omega=2\pi[1-\cos(\Delta\theta^0)]=2\pi(1-\hat{n}_0\cdot\hat{r}_0)=2\pi\left(\frac{1}{10}-\frac{1}{15}\right). \qquad (151)$$

This background reduction may suffice for facilitating a R^0 research.

If the hadronic showers in the special events which do correlate with decays in the detector D are indeed due to R^0 decays they should extend further into the calorimeters than the average hadronic jets. If this feature which we expect on the basis of our discussion of the R^0 evolution is confirmed it will provide further evidence for R^0 particles.

Needless to say all these very qualitative arguments need verification via detailed Monte Carlo studies with $\tilde{g}$ jets and evolving R^0 particles. The whole analysis can be repeated for e^+e^- collisions at the Z^0 resonance. We need now to specifically look for 4 jet events: $e^+e^-\rightarrow q\bar{q}\tilde{g}\bar{\tilde{g}}$ which may be relatively rare occurring in only few percent of the events. However we gain by having a much reduced hadronic background and a larger δt. [The very large e.m. radiative background can be effectively avoided by thick lead shieldings around the detector.]

In summary we have suggested a novel, and in our view optimized, alternative R^0 search method. Instead of searching for R^0 decays in fixed target neutral beam where the putative R^0's are at a considerable disadvantage, we suggest using colliders and jets at 90° where R^0 could be relatively copious. Also by moderating first the R^0 particles a very wide range of R^0 lifetimes becomes accessible.

While many of the basic R^0 features appear in the papers by Farrar we included them for completeness. The discussion of Secs. IV and V have new elements, and in particular the possibility of having longer attenuation path for R^0 in matter as compared with neutrons is discussed at length.

Negative results in the experiments suggested cannot unequivocally exclude the eight gluino hypothesis. Though we argued for R^0N cross sections smaller than N-N cross sections the reverse may be true and the R^0 will then migrate less than our estimates making it more difficult to separate R^0 decays from the ''normal'' hadronic activity. Also for $\tau_{R^0}\simeq3\times10^{-9}-3\times10^{-8}$ sec direct searches for decay in flight may be more efficient.

Eventually many different lines of research will converge to yield a definitive verdict on the light gluino hypothesis. At present, however, the R^0 may exist. It will be a great pity if it will be missed due to lack of enthusiasm and possibly for not trying the right optimal approach to discovering it. After completing this work a preprint by R. Mohapatra and S. Nandi (UMD-PP-97-082) addressing light gluinos came to our attention.

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