Fermi National Accelerator Laboratory

FERMILAB-PUB-90/105-T
May 1990

Cosmions with spin–dependent interactions

Gian F. Giudice
Fermi National Accelerator Laboratory
P.O. Box 500, Batavia, IL 60510
and
Stuart Raby
Department of Physics, The Ohio State University
Columbus, OH 43210

Abstract

If the cosmion, the particle proposed to solve both the solar neutrino and the dark matter problem, has a purely spin–dependent interaction with nuclei, it would remain invisible to presently existing halo particle detectors. We consider the ingredients necessary for building an acceptable model for a spin–dependent interacting cosmion, and propose a concrete example. This particular model introduces a new neutral gauge boson, lighter than the Z, which is shown to be consistent with low–energy neutral current data and could be lurking in presently operating colliders.

Operated by Universities Research Association Inc. under contract with the United States Department of Energy

I.

Faulkner, Gilliland, Press, and Spergel [1] have shown that the observed deficit [2] in the solar neutrino flux can be explained by postulating the presence of a new particle in the interior of the Sun. The effect of this weakly interacting particle – generally referred to as the cosmion – is to transport heat away from the solar core, and thus to slightly lower its temperature. The solar neutrino flux, which strongly depends on the core temperature, gets reduced with respect to the standard prediction, in agreement with observations.

In order to provide the correct depletion of the neutrino flux, the cosmion should have an average cross section per baryon in the Sun in the range $\sigma = 10^{-36}$–10^{-33} cm^2 [3] and a mass between 2–4 GeV* and 8 GeV. A sufficient solar concentration of cosmions can be achieved if the cosmion density in the neighborhood of the Sun is of the order of the halo density. This interesting coincidence suggests a connection between the solar neutrino and dark matter problem. Such an identification further requires that the cosmion is stable (in cosmological times) with a highly suppressed annihilation cross section inside the Sun, $\sigma_A \lesssim 10^{-4}\sigma$.

The requirements defined above (especially the need for a σ somewhat larger than the typical weak cross section) tightly constrain the cosmion properties, and no known elementary particle is a suitable candidate for it. Nevertheless, several authors [5] have suggested different extensions of the standard model in which a particle with the desired properties can be found.

Caldwell et al. [6] have recently improved the limits on searches for halo particles by means of a Silicon detector, and have set an upper bound of about 4–6 GeV (depending on σ) on the mass of the cosmion. Although the cosmion is not completely ruled out, the acceptable window in the mass range has been severely narrowed. However, it is important to bear in mind that the experiment of ref.[6] is sensitive only to particles with spin–independent interactions; in this paper, we want to propose a new class of models, where the cosmions have purely spin–dependent interactions with nuclei, thus evading the present experimental limits.

II.

The most familiar example of a particle with a nuclear cross section that is purely spin-dependent is a Majorana fermion (X), with an effective interaction with ordinary quarks

*The lower bound on the cosmion mass, which corresponds to the evaporation mass in the Sun, is a function of the cosmion–nucleus cross section [4]. For the values of σ considered in the following ($\sigma \simeq 10^{-36}$ cm^2), the limit on the cosmion mass is about 4 GeV.

and leptons (f) of the form

$$\mathcal{L}_I = \sum_f \frac{1}{\Lambda_f^2} \bar{X}\gamma^\mu\gamma_5 X \bar{f}\gamma_\mu(g_f + \gamma_5)f \quad . \tag{1}$$

In eq.(1), Λ_f is a dimensionful parameter that measures the energy scale of the interaction and g_f is a dimensionless coupling constant. The scattering cross section of the particle X in the Sun is dominated by its interaction with hydrogen. In the non-relativistic limit, the X-proton cross section is given by:

$$\sigma = \frac{12\ m_p^2 m_X^2}{\pi(m_p + m_X)^2}\left(\sum_q \frac{1}{\Lambda_q^2}\Delta q\right)^2 \quad , \tag{2}$$

where the sum is over quarks; m_p and m_X are the proton and X mass, and Δq is the component of the proton spin carried by quarks. Using the results by the EMC collaboration [7], together with data from neutron lifetime and weak decays of hyperons, one finds:

$$\Delta u = 0.75 \pm 0.06 \quad \Delta d = -0.50 \pm 0.06 \quad \Delta s = -0.22 \pm 0.06 \quad . \tag{3}$$

For X to be a possible cosmion candidate, the cross section σ in eq.(2) must satisfy

$$X_H \sigma \gtrsim 10^{-36}\mathrm{cm}^2 \quad , \tag{4}$$

where X_H is the hydrogen fraction of the solar mass (77%). Assuming that the up quark has the largest coupling in the four-fermion interaction of eq.(1), eq.(4) then yields:

$$\Lambda_u \lesssim 150\ \mathrm{GeV}\left(1 + \frac{m_p}{m_X}\right)^{-1/2} \quad . \tag{5}$$

Before interpreting X as a cosmion, we have to make sure that it does not annihilate too rapidly either in the early Universe, thus depleting the relic abundance, or in the interior of the Sun, thereby preventing a solution of the solar neutrino problem. The interaction (1) leads to the annihilation process $XX \to f\bar{f}$, where the non-relativistic cross section is

$$\sigma_A(XX \to f\bar{f})v = \frac{2}{\pi}m_X^2 \sum_f \frac{c_f}{\Lambda_f^4}\sqrt{1-y}\left\{y + v^2\left[\frac{1+g_f^2}{3} - y\frac{7-2g_f^2}{12} + \frac{y^2}{8(1-y)}\right]\right\} \quad . \tag{6}$$

In eq.(6), c_f is the color factor, v is the relative velocity of the annihilating particles (in units of c) and $y \equiv m_f^2/m_X^2$.

In the case of annihilation inside the Sun, v is given by the average velocity of halo particles ($\simeq 10^{-3}$), and the rate in eq.(6) is dominated by the "s-wave" term ($v = 0$). Requiring that, inside the Sun, the annihilation cross section is strongly suppressed with respect to the scattering cross section, $\sigma_A(XX \to f\bar{f})v \lesssim 10^{-40}\text{cm}^2$, we obtain:

$$\Lambda_f \gtrsim 1.2 \text{ TeV} \left(c_f \sqrt{1 - \frac{m_f^2}{m_X^2}} \right)^{1/4} \left(\frac{m_f}{\text{GeV}} \right)^{1/2} . \tag{7}$$

Therefore, eq.(7) implies that charged leptons and quarks of the second and third generation (if lighter than X) should necessarily have an effective interaction with X much weaker than the one required in eq.(5). Moreover, a limit on the effective interaction between X and the electron can be extracted from data on single-photon production in e^+e^- collisions. Using the analysis of ASP [8], we derive the limit $\Lambda_e > 270$ GeV for $m_X = 0$ (or $\Lambda_e > 190$ GeV for $m_X = 10$ GeV).

We are therefore led to believe that the dominant X annihilation mode in the early Universe (where the "p-wave" terms become important) is into either light quarks or neutrinos. An upper bound on the X relic abundance (in units of the critical density), Ω_X, can be derived by considering just the annihilation into up quarks. We estimate:

$$\Omega_X \simeq \frac{2 \cdot 10^{-37}\text{cm}^2}{h^2\, \sigma_A v} \simeq 0.2 \left(\frac{\Lambda_u}{150 \text{ GeV}} \right)^4 \left(\frac{5 \text{ GeV}}{m_X} \right)^2 \left(\frac{1/2}{h} \right)^2 , \tag{8}$$

where h is the Hubble constant in units of 100 km s^{-1} Mpc^{-1} and $\sigma_A v$ is evaluated at the freeze-out temperature, where $v^2 \simeq 3/10$. Eq.(8) shows that the particle X can be an acceptable dark matter candidate, even if the energy scale Λ of the interaction (1) is as low as the one required by eq.(5), provided that the annihilation into fermion species different from light quarks is suppressed.

In summary, a Majorana fermion X with an effective coupling with quarks and leptons given by eq.(1) can be a possible cosmion candidate with a purely spin-dependent interaction with nuclei if: *i)* the interaction with either up or down quarks has a strength of order $\Lambda \simeq 150$ GeV, and *ii)* the interactions with charged leptons and heavy quarks are suppressed, typically by one order of magnitude (see eq.(7)).

An effective interaction of the form (1) can be produced by the exchange of new scalar particles or of new vector bosons. These scalar particles, necessarily colored, would produce large missing p_T signals in high energy hadron colliders [9] and can be ruled out by the CDF data from Tevatron. A new light gauge boson Z' can be phenomenologically acceptable, provided its couplings to leptons are suppressed, as discussed in the next section. However, even if one can understand the strength of the new interaction in terms of some exotic particles, the different couplings to the various quark generations cause a serious problem. Breaking flavor symmetry at low energy naturally leads to unacceptable flavor changing neutral currents, which can be avoided only at the price of some peculiar fine tunings among parameters. Therefore, we abandon the idea of a Majorana cosmion and we turn to discuss the possibility of a Dirac particle.

III.

If the cosmion is a Dirac fermion, the bounds from annihilation can be easily relaxed by assuming a cosmic asymmetry (ϵ_X) of the same order of the baryon asymmetry. Now the relic abundance is fixed by the value of ϵ_X, and the ratio between luminous and dark matter is naturally explained, and is of order m_p/m_X. Moreover, nearly no anticosmions are left at present, and thus no cosmion annihilation in the Sun can occur.

In order to obtain a large spin-dependent cosmion–nucleus scattering cross section, we still need an effective interaction of the form (1). We will therefore assume the existence of a new neutral gauge boson, typically lighter than the Z. Such a Z' must satisfy the following requirements: *i)* it should have vanishing couplings to neutrinos and charged leptons, in order to be consistent with the precise measurements on low energy neutral current processes, *ii)* it should have only an axial–vector coupling to either the cosmion or to all the quarks, to provide the purely spin–dependent X-nucleus interaction, and *iii)* it should have a small enough mixing with the ordinary Z, or else it modifies the observed properties of the Z gauge boson.

Let us now describe a specific model with an extra $U(1)$ gauge group which satisfies all these requirements; we start by enlarging the standard model to include a new weak singlet Dirac fermion X, the cosmion. We then choose an extended Higgs sector, containing three weak doublets and one singlet, where the Yukawa interactions have the form:

$$\mathcal{L}_Y = \lambda_u \bar{u}_R q_L \Phi_u + \lambda_d \bar{d}_R q_L \Phi_d + \lambda_e \bar{e}_R l_L \Phi_e + \lambda_X \bar{X}_R X_L \Phi_X + h.c. \ . \tag{9}$$

In eq.(9), summation over generation indices is understood. The quantum numbers Y' of the new gauge interaction for ordinary left– and right–handed quarks (q_L, u_R, d_R) and leptons (l_L, e_R), for the cosmion (X), and for the three weak doublets (Φ_u, Φ_d, Φ_e) and the one singlet (Φ_X) of Higgs fields, are shown in Table 1. The values of K_q, K_1, K_2, and K_X in Table 1 are left arbitrary. Note that gauge invariance (for $K_1 + K_2 \neq 0$) insures that eq.(9) is the most general Yukawa interaction and it prevents the occurrence of tree level flavor changing neutral currents due to Higgs exchange.

The mass terms for the gauge bosons are:

$$\mathcal{L}_m = \sum_{\Phi} < \Phi^\dagger > \left| \frac{g}{2}\vec{\tau} \cdot \vec{W}_\mu + \frac{g'}{2} Y B_\mu + \frac{g''}{2} Y' Z'_\mu \right|^2 < \Phi > \ , \tag{10}$$

where g'' is the coupling constant of the new $U(1)$ gauge group and the sum in eq.(10) runs over the different Higgs fields Φ. The vacuum expectation values $< \Phi >$ which preserve electric charge ($Q = I_{(3)} + Y/2$) are:

$$< \Phi_u > \equiv \begin{pmatrix} 0 \\ v_u \end{pmatrix} \quad < \Phi_d > \equiv \begin{pmatrix} v_d \\ 0 \end{pmatrix} \quad < \Phi_e > \equiv \begin{pmatrix} v_e \\ 0 \end{pmatrix} \quad < \Phi_u > \equiv v_X \ . \tag{11}$$

By substituting eq.(11) into eq.(10), we obtain a charged W boson with mass $m_W^2 = g^2(v_u^2 + v_d^2 + v_e^2)/2$, a massless photon $A \equiv \sin\theta_W W^{(3)} + \cos\theta_W B$, $\tan\theta_W \equiv g'/g$, and the following mass term for the $Z \equiv -\cos\theta_W W^{(3)} + \sin\theta_W B$ and Z' gauge bosons:

$$\mathcal{L}_m = \frac{1}{2}\frac{m_W^2}{\cos^2\theta_W} \begin{pmatrix} Z^\mu & Z^{\mu\prime} \end{pmatrix} \begin{pmatrix} 1 & \delta \\ \delta & \beta \end{pmatrix} \begin{pmatrix} Z_\mu \\ Z'_\mu \end{pmatrix} \ , \tag{12}$$

$$\beta \equiv \frac{g''^2}{g^2 + g'^2} \frac{K_1^2 v_u^2 + K_2^2 v_d^2 + 4K_X^2 v_X^2}{v_u^2 + v_d^2 + v_e^2} \qquad \delta \equiv \frac{g''}{\sqrt{g^2 + g'^2}} \frac{K_1 v_u^2 - K_2 v_d^2}{v_u^2 + v_d^2 + v_e^2} \ . \tag{13}$$

Notice that the mixing term δ between Z and Z' turns out to be small if the minimization of the scalar potential yields vacuum expectation values such that:

$$\frac{v_u}{v_d} \simeq \sqrt{\frac{K_2}{K_1}} \ . \tag{14}$$

As we will show in the following, the relation (14) does not need to be satisfied to high accuracy (see eq.(22)) to make the model consistent with the measurements of weak processes. Therefore, even though we do not attempt to write down the complete scalar potential of the theory, we believe that eq.(14) should not necessarily be regarded as an unnatural fine tuning of our model.

In the limit of small δ, the non-relativistic cross section for cosmion-proton scattering is:

$$\sigma = \frac{3g''^4}{64\pi m_{Z'}^4} \frac{m_p^2 m_X^2}{(m_p + m_X)^2} K_X^2 \left[K_1 \,\Delta u + K_2(\Delta d + \Delta s)\right]^2 \ . \tag{15}$$

Now the constraint of eq.(4) implies that the mass of the Z' boson must be:

$$m_{Z'} \lesssim 30 \text{ GeV } \frac{g''}{g} \sqrt{\frac{K_X\left[K_1\,\Delta u + K_2(\Delta d + \Delta s)\right]}{1 + \frac{m_p}{m_X}}} \ . \tag{16}$$

Our model can account for such a light Z' by having v_e (which gives masses to Z and W, but not to Z') much larger than the other vacuum expectation values in eq.(11).

At this stage, the model has local gauge anomalies because of the $U(1)'$ current, as is apparent from the Y' quantum number assignment in Table 1; we are therefore forced to introduce new fermions in order to cancel the anomalies. Embedding the model in some grand unified scheme is a natural way to implement the anomaly cancellation. For instance, let us consider the popular 27-dimensional representation of the anomaly-free group $E(6)$, which contains one family of quarks and leptons, a weak-singlet Dirac particle (interpreted as the cosmion), two weak-singlet color-triplet and two weak-doublet color-singlet Weyl fermions. There are two possible ways to choose a low energy extra $U(1)$ with vanishing quantum numbers for charged leptons and neutrinos, corresponding to the two ways of assigning hypercharge to the model. The first option is the choice proposed by Ross and Segrè[†] [5], which corresponds to

$$K_1 = K_X = 1 \quad , \quad K_2 = 0 \quad , \quad K_q = -1/3 \quad , \tag{17}$$

up to an overall normalization absorbed in g''. The second option, obtained by flipping the assignment of the right-handed fields, does not fall in our class of models, since it corresponds to a Z' interaction which is not purely axial-vector with either X or with all quarks.

In the case of eq.(17), since $K_2 = 0$, we are no longer able to set $\delta = 0$ and keep non-zero vacuum expectation values, as in eq.(14). Nevertheless, we will show in the following that it is still possible to have a Z–Z' mixing small enough to be consistent with neutral current data, if $m_{Z'}$ is not too large.

A potential problem of our model may arise because the new fermions required by anomaly cancellation can not be made arbitrarily heavy. In fact their masses are protected by the $U(1)'$ symmetry which is unbroken up to very low energy scales. Eq.(16) implies a bound on v_u and v_d which, assuming all Yukawa couplings $\lambda < \sqrt{4\pi}$, leads to the constraint

[†] Ross and Segrè have considered only a cosmion with coherent interactions with nuclei. However, if the fields ν_4 and ν_5, in their notation, are interpreted as a Dirac cosmion, and the Y' quantum numbers are as in eq.(17), then their model coincides with ours.

that all new fermions required by anomaly cancellation should have a mass $m < 180$ GeV. This completes the description of the model. Now we want to study the effect of the new light Z' in measurements of neutral current weak interactions and in searches for exotic signals in high-energy collisions.

First, let us see how large a deviation from eq.(14) can be tolerated. Even if the Z' is not coupled to leptons, its mixing with the Z will modify the standard predictions for leptonic and semileptonic processes. If $\delta \neq 0$, the weak current coupled to the Z boson is:

$$J_Z^\mu = \frac{g}{2\cos\theta_W}\bar{f}\gamma^\mu(c_V - c_A\gamma_5)f \tag{18}$$

$$c_V = (I_{(3)} - 2\sin^2\theta_W Q)\cos\tilde{\theta} + \frac{g''}{g}\cos\theta_W\left(\frac{Y_L' + Y_R'}{2}\right)\sin\tilde{\theta} \tag{19}$$

$$c_A = I_{(3)}\cos\tilde{\theta} + \frac{g''}{g}\cos\theta_W\left(\frac{Y_L' - Y_R'}{2}\right)\sin\tilde{\theta} \tag{20}$$

$$\tan 2\tilde{\theta} = \frac{2\delta}{\beta - 1}\ , \tag{21}$$

where $I_{(3)}$ is the third component of isospin and Q is the charge of the fermion f, and $Y_{L,R}'$ are the $U(1)'$ quantum numbers for the left- and right-handed components ($Y_L' = K_q$, $Y_R' = K_q + K_1$ for up quarks, $Y_L' = K_q$, $Y_R' = K_q + K_2$ for down quarks, $Y_L' = Y_R' = 0$ for leptons). The ρ-parameter, measuring the ratio between the strength of neutral and charged current interaction, is also modified by the Z–Z' mixing:

$$\rho - 1 \equiv \frac{m_W^2}{\cos^2\theta_W m_Z^2} - 1 = \frac{m_Z^2}{m_{Z'}^2 - m_Z^2}\delta^2 + O(\delta^4)\ . \tag{22}$$

In terms of the mixing angle $\tilde{\theta}$ and the parameter ρ, one can compute all the deviations from the standard model predictions on neutral current data[†]. Table 2 contains the expressions, in leading order in δ, for the extra contributions to the standard model parameters measured in deep-inelastic ν scattering, $\nu_\mu e$ scattering, atomic parity and polarized eD asymmetry experiments, and to the Z properties measured by the LEP detectors. The table shows that, for $\frac{g''}{g}\delta \lesssim 0.06$, the existence of a Z' as light as in eq.(16) is consistent with the present experimental measurements, well within a 2σ experimental error. In turn, this limit

[†] Here we are neglecting contributions from radiative corrections, and concentrate only on the effect of Z–Z' mixing.

on δ implies:

$$\frac{|K_1 v_u^2 - K_2 v_d^2|}{K_1^2 v_u^2 + K_2^2 v_d^2 + 4K_X^2 v_X^2} \simeq \frac{g'' \cos\theta_W}{g} \frac{m_Z^2}{m_{Z'}^2} \delta \lesssim \frac{1}{2} \left(\frac{m_Z}{3m_{Z'}}\right)^2 \quad , \tag{23}$$

The inequality (23) shows that, for a light Z', the model is consistent with the measurement on weak processes, even if eq.(14) is verified without a high degree of accuracy. Larger values of $m_{Z'}$ require that condition (14) should be satisfied with a better accuracy. Furthermore, the case of the Ross-Segrè model, eq.(17), is also acceptable, if v_X is not smaller than v_u and $m_{Z'}$ is not too large.

Next, we consider the detection of Z' bosons in high-energy colliders. The Z' can decay into a quark-antiquark pair with a branching ratio $B_h = f/(1+f)$, or invisibly into cosmions, giving rise to a missing energy signal, with a branching ratio $B_c = 1/(1+f)$, where

$$f = \frac{3}{2K_X^2} \left[5K_q^2 + 2(K_q + K_1)^2 + 3(K_q + K_2)^2\right] \quad . \tag{24}$$

In hadronic machines, the Z' can be produced in a Drell-Yan process. The quark-antiquark decay mode can be detected as a bump over the QCD background in the di-jet invariant mass distribution. The UA1 collaboration has published a limit on the product of the cross section with the decay branching ratio ($\sigma_P B_h$) for production of a particle with mass m decaying into two hadronic jets, in the region $m > 150$ GeV [11]. In absence of a limit for $m < 150$ GeV, the authors of ref.[12] have suggested a simple estimate for a bound on $\sigma_P B_h$ by using UA2 data [13] for di–jet production cross section. For a particle with width $\Gamma < 0.1\ m$, they estimate $\sigma_P B_h \lesssim 20$ nb for $m = 50$ GeV and $\sigma_P B_h \lesssim 10$ nb for $m = 80$ GeV (at $\sqrt{s} = 630$ GeV). With the choice of parameters of eq.(17), the cross section for Z' production at $\sqrt{s} = 630$ GeV is $\sigma_P B_h = g''^2/g^2 \cdot 3.9$ nb for $m = 50$ GeV and $\sigma_P B_h = g''^2/g^2 \cdot 0.9$ nb for $m = 80$ GeV. For $\Gamma > 0.1\ m$ (*i.e.* $g''/g > 3.1$, in the case of eq.(17)) the resonance is broader than the experimental energy resolution and separation between signal and background seems much harder. No limit can be derived for $m < 50$ GeV.

The decay of the Z' into cosmions will provide an excess of missing transverse energy in hadron collisions. The UA1 collaboration, studying monojet events with significant missing transverse energy, has published a limit [14] on the number of neutrino families, $N_\nu < 10$ at 90% C.L. Even if the light Z' will produce a missing energy spectrum softer than the one from $Z \to \bar{\nu}\nu$, the UA1 limit can constrain a model for which both g''/g and B_c are large. In addition the CDF collaboration is now performing an analysis on invisible Z decays, which could provide a limit on the extra Z'.

A light Z' could also be produced at LEP in the Z decay modes $Z \to Z'\bar{q}q$ or $Z \to Z'\gamma$. In the limit of massless quarks, the decay width for $Z \to Z'\bar{q}q$ is:

$$\Gamma(Z \to Z'\bar{q}q) = \frac{4\pi^2\alpha^2}{\sin^4\theta_W \cos^2\theta_W m_Z}\left(\frac{g''}{g}\right)^2 \cdot$$

$$\cdot \left[2K_q^2(c_A^2 + c_V^2) + K_i(K_i + 2K_q)(c_A - c_V)^2\right] \int d\Phi^{(3)} M^2 \ , \tag{25}$$

$$M^2 = \frac{1}{2}(d_1 d_2 - \beta)\left(\frac{1}{d_1^2} + \frac{1}{d_2^2}\right) + \frac{1+\beta}{d_1 d_2}(1 + \beta - d_1 - d_2) \ . \tag{26}$$

Here c_A and c_V are the quark couplings to the Z, given in eqs.(19–20) for $\delta = 0$; K_i is equal to K_1 (K_2) for up- (down-) type quarks; $\beta = m_{Z'}^2/m_Z^2$; $d_a = 1 - 2E_a/M_Z$, and E_a ($a = 1,2$) are the energies of quark and antiquark in the Z rest frame; $d\phi^{(3)}$ is the three-body phase space measure. When the Z' decays into cosmions, the process gives the identifiable signature $Z \to$ jets+ $\not{E}_T$. In fig.1, we show the product $\mathrm{BR}(Z \to Z'\bar{q}q)\cdot\mathrm{BR}(Z' \to \bar{X}X)$, summed over all quarks except the top where, for illustrative purposes, the choice of Y' quantum numbers of eq.(17) is assumed. With cuts on $\not{E}_T$ and jet acoplanarity, the standard model background can be disentangled, and the signal can be observed at LEP down to Z branching ratios of order 10^{-6}. Recently the ALEPH collaboration has reported a limit of about 10^{-4} on the Z branching ratio into jets and large missing transverse energy [15].

The process $Z \to Z'\gamma$ is mediated by loops of quarks and of the new fermions which are necessary to cancel the anomalies. Summing over intermediate fermions (f) with electric charge Q, color factor c_f and mass m_f, the partial width is:

$$\Gamma(Z \to Z'\gamma) = \frac{\alpha^3 m_Z}{96\pi^2 \sin^4\theta_W \cos^2\theta_W}\left(\frac{g''}{g}\right)^2 (1-\beta)^3\left(1 + \frac{1}{\beta}\right) \cdot$$

$$\cdot \left\{\sum_f c_f Q\left[c_A(2K_q + K_i) h\left(\frac{m_Z^2}{m_{Z'}^2}, \frac{m_f^2}{m_{Z'}^2}\right) + c_V K_i\left(h\left(\frac{m_{Z'}^2}{m_Z^2}, \frac{m_f^2}{m_Z^2}\right) - \frac{1}{2}\right)\right]\right\}^2 \ , \tag{27}$$

$$h(a,b) = \int_0^1 dx \int_0^{1-x} dy \left[1 + a\frac{1-x-y}{x} - \frac{b}{xy}\right]^{-1} \quad \text{and} \tag{28}$$

$$h(a,b) \simeq \frac{-1}{24b} \;\; for \;\; b >> a,1 \;\; h(a,b) \simeq \frac{1}{2(1-a)}\left(1+\frac{a}{1-a}\log a\right) \;\; for \;\; b << a,1. \qquad (29)$$

The Z' decay into cosmions generates a highly characteristic event, free of standard model background, with a single photon of energy $E_\gamma = m_Z(1-\beta)/2$ recoiling against missing energy. The contribution of light quarks in eq.(27) is suppressed in the limit $m_{Z'} << m_Z$ and the exchange of heavy fermions, which do not decouple, will typically give rates too small to be observed at LEP.

Fig.2 summarizes the region of parameters g'' and $m_{Z'}$ leading to a solution for the cosmion, which can be studied at LEP via the processes $Z \rightarrow$ jets+ $\not{E}_T$. Part of the allowed region of parameters could also be excluded by UA1 and UA2 extending to a lower energy range their analysis for di-jet mass distribution and for missing transverse momentum events, as we have previously discussed.

IV.

In conclusion, we have considered the possibility of elementary particle models with a cosmion candidate having only a spin-dependent interaction with nuclei, and thus evading the present experimental search. Since the cosmion should interact with matter via an axial-vector current, a Majorana fermion is the most natural candidate. However, the constraints coming from the annihilation in the Sun and in the early Universe require that the cosmion should have an effective interaction with light quarks very different from the one to heavy quarks. Therefore, the cosmion-quark current interaction will generally lead to flavor changing neutral currents, making the particle models unacceptable.

On the other hand, if the cosmion is a Dirac particle, the limits from annihilation can be evaded by assuming a cosmion-anticosmion asymmetry. We have proposed a model for a Dirac cosmion, which has a purely spin-dependent interaction with nuclei through a new light neutral gauge boson. A particular case of our model corresponds to an early proposal by Ross and Segrè [5], in the context of cosmions with spin-independent interactions. We have also found a mechanism to prevent large mixings between the Z and Z', and have shown that the existence of such a light Z' is consistent with neutral current data. The major difficulty with the model is the need for new heavy fermions to cancel the anomaly in the Z' current. An optimistic upper limit on the masses of these new fermions (including the top quark) is about 180 GeV; otherwise Yukawa interactions become non-perturbative. Signals for the new Z' can be discovered at high-energy colliders and, in particular, LEP can severely test our model. Small corners of parameter space, however, can survive searches at

LEP and other colliders, thus the direct search for spin-dependent interacting halo particles remains a goal of primary importance. In conclusion, the cosmion is still a viable solution to the solar neutrino and dark matter problem.

Acknowledgements

G.F.G. wants to thank J. Frieman for helpful discussions and the Department of Physics of The Ohio State University for kind hospitality.

References

[1] D.N. Spergel and W.M. Press, Ap. J.**294** (1985) 663; J. Faulkner and R.L. Gilliland, Ap. J. **233** (1985) 994; R.L. Gilliland, J. Faulkner, W.M. Press and D.N. Spergel, Ap. J. **306** (1986) 703.

[2] R. Davis, *et al.*, in Proceedings of the 21st International Cosmic Ray Conference, Adelaide, Australia (1990); K.S. Hirata, *et al.*, Phys. Rev. Lett. **63** (1989) 16.

[3] A. Gould and G. Raffelt, Princeton Preprint AST 89/20 (1989).

[4] A. Gould, Princeton Preprint AST 89/49 (1989).

[5] G.B. Gelmini, L.J. Hall and M.J. Lin, Nucl. Phys. **B281** (1987) 726; S. Raby and G.B. West, Nucl. Phys. **B292** (1987) 793; S. Raby and G.B. West, Phys. Lett. **194B** (1987) 557; G.G. Ross and G. Segrè, Phys. Lett. **197B** (1987) 45; S. Raby and G.B. West, Phys. Lett. **202B** (1988) 47; G.F. Giudice and E. Roulet, Phys. Lett. **219B** (1989) 309.

[6] D. Caldwell, *et al.*, in preparation (1990).

[7] J. Ashman, *et al.* (EMC collaboration), Nucl. Phys. **B328** (1989) 1.

[8] C. Hearty, *et al.*, Phys. Rev. **D39** (1989) 3207.

[9] G.F. Giudice, G. Ridolfi and E. Roulet, Phys. Lett. **211B** (1988) 370.

[10] Particle Data Group, G.P. Yost, *et al.*, Phys. Lett. **204B** (1988) 1.

[11] C. Albajar, *et al.* (UA1 collaboration), Phys. Lett. **209B** (1988) 127.

[12] S. Dimopoulos, *et al.*, Phys. Rev. **D41** (1990) 2099.

[13] R. Ansari, *et al.* (UA2 collaboration), Phys. Lett. **186B** (1987) 452.

[14] C. Albajar, *et al.* (UA1 collaboration), Phys. Lett. **185B** (1987) 241.

[15] D. Decamp, *et al.* (ALEPH collaboration), Preprint LAL 90-26, CERN-EP/90-63 (1990).

Table Captions

Table 1. Y and Y' quantum number assignment for left- and right-handed quarks, leptons and cosmion, and for the Higgs fields.

Table 2. The table contains the expressions for the extra contributions from Z-Z' mixing, at leading order in δ, to the standard model parameters, as defined by the Particle Data Group [10], to the hadronic (Γ_h) and electronic (Γ_e) Z decay width and to the effective number of neutrino families (N_ν) measured from the shape of the Z resonance. We use the notation $s_W \equiv \sin\theta_W$, $c_W \equiv \cos\theta_W$ and $\beta \equiv m_{Z'}^2/m_Z^2$. We also give the numerical values for these extra contributions, taking $g''/g = 1$, $\sin^2\theta_W = 0.23$, $m_Z = 90.1$ GeV and $\beta = 1/9$. Finally, in the last column, we give the present experimental error on the standard model parameters, at the level of 2σ [10]. For Γ_h, Γ_e and N_ν, we give the expected sensitivity of the LEP experiments, again at the level of 2σ.

Figure Captions

Fig.1. $(g/g'')^2 BR(Z \to Z'\bar{q}q) \cdot BR(Z' \to \bar{X}X)$ for $K_1 = K_X = 1$, $K_2 = 0$, $K_q = -1/3$.

Fig.2. Values of $BR_q \equiv BR(Z \to Z'\bar{q}q) \cdot BR(Z' \to \bar{X}X)$ in the $g''/g - m_{Z'}$ region allowed by the cosmion solution, for $K_1 = K_X = 1$, $K_2 = 0$, $K_q = -1/3$.

Table 1:

	q_L	u_R	d_R	l_L	e_R	X_L	X_R	Φ_u	Φ_d	Φ_e	Φ_X
Y	1/3	4/3	-2/3	-1	-2	0	0	1	-1	-1	0
Y'	K_q	$K_q + K_1$	$K_q + K_2$	0	0	$-K_X$	K_X	K_1	K_2	0	$2K_X$

Table 2:

	extra contribution from Z-Z' mixing	numerical value	2 σ exp. error
$\epsilon_L(u)$	$-\frac{g''}{g}\frac{c_W}{2}K_q\frac{\delta}{1-\beta}$	$-0.49\ K_q\delta$	0.034
$\epsilon_L(d)$	$-\frac{g''}{g}\frac{c_W}{2}K_q\frac{\delta}{1-\beta}$	$-0.49\ K_q\delta$	0.028
$\epsilon_R(u)$	$-\frac{g''}{g}\frac{c_W}{2}(K_1+K_q)\frac{\delta}{1-\beta}$	$-0.49\ (K_1+K_q)\delta$	0.028
$\epsilon_R(d)$	$-\frac{g''}{g}\frac{c_W}{2}(K_2+K_q)\frac{\delta}{1-\beta}$	$-0.49\ (K_2+K_q)\delta$	0.138
g_L^2	$\frac{g''}{g}\frac{c_W}{3}s_W^2K_q\frac{\delta}{1-\beta}$	$0.07\ K_q\delta$	0.009
g_R^2	$\frac{g''}{g}\frac{c_W}{3}s_W^2(2K_1-K_2+K_q)\frac{\delta}{1-\beta}$	$0.07\ (2K_1-K_2+K_q)\delta$	0.007
g_V^e	$(\frac{1}{2}-2s_W^2)\frac{(2-\beta)}{(1-\beta)^2}\delta^2$	$0.09\ \delta^2$	0.072
g_A^e	$\frac{(2-\beta)}{2(1-\beta)^2}\delta^2$	$1.19\ \delta^2$	0.054
C_{1u}	$\frac{g''}{g}\frac{c_W}{2}(K_1+2K_q)\frac{\delta}{1-\beta}$	$0.49\ (K_1+2K_q)\delta$	0.142
C_{1d}	$\frac{g''}{g}\frac{c_W}{2}(K_2+2K_q)\frac{\delta}{1-\beta}$	$0.49\ (K_2+2K_q)\delta$	0.128
$C_{2u}-\frac{1}{2}C_{2d}$	$-\frac{g''}{g}\frac{c_W}{2}(1-4s_W^2)(K_1-\frac{1}{2}K_2)\frac{\delta}{1-\beta}$	$-0.02\ (2K_1-K_2)\delta$	0.74
Γ_h	$\frac{G_F m_Z^3}{\sqrt{2}\pi}\frac{g''}{g}c_W[(\frac{1}{2}+\frac{2}{3}s_W^2)K_q+$ $+s_W^2(\frac{4}{3}K_1-K_2)]\frac{\delta}{1-\beta}$	$436\ (1.3K_1-K_2+$ $+2.8K_q)\delta$ MeV	40 MeV
Γ_e	$-\frac{G_F m_Z^3}{12\sqrt{2}\pi}(1-4s_W^2+8s_W^4)\frac{(2-\beta)}{(1-\beta)^2}\delta^2$	$-192\ \delta^2$ MeV	2 MeV
N_ν	$[2(\frac{g''}{g})^2c_W^2K_X^2-3(2-\beta)](\frac{\delta}{1-\beta})^2$	$(1.9K_X^2-7.2)\delta^2$	0.1

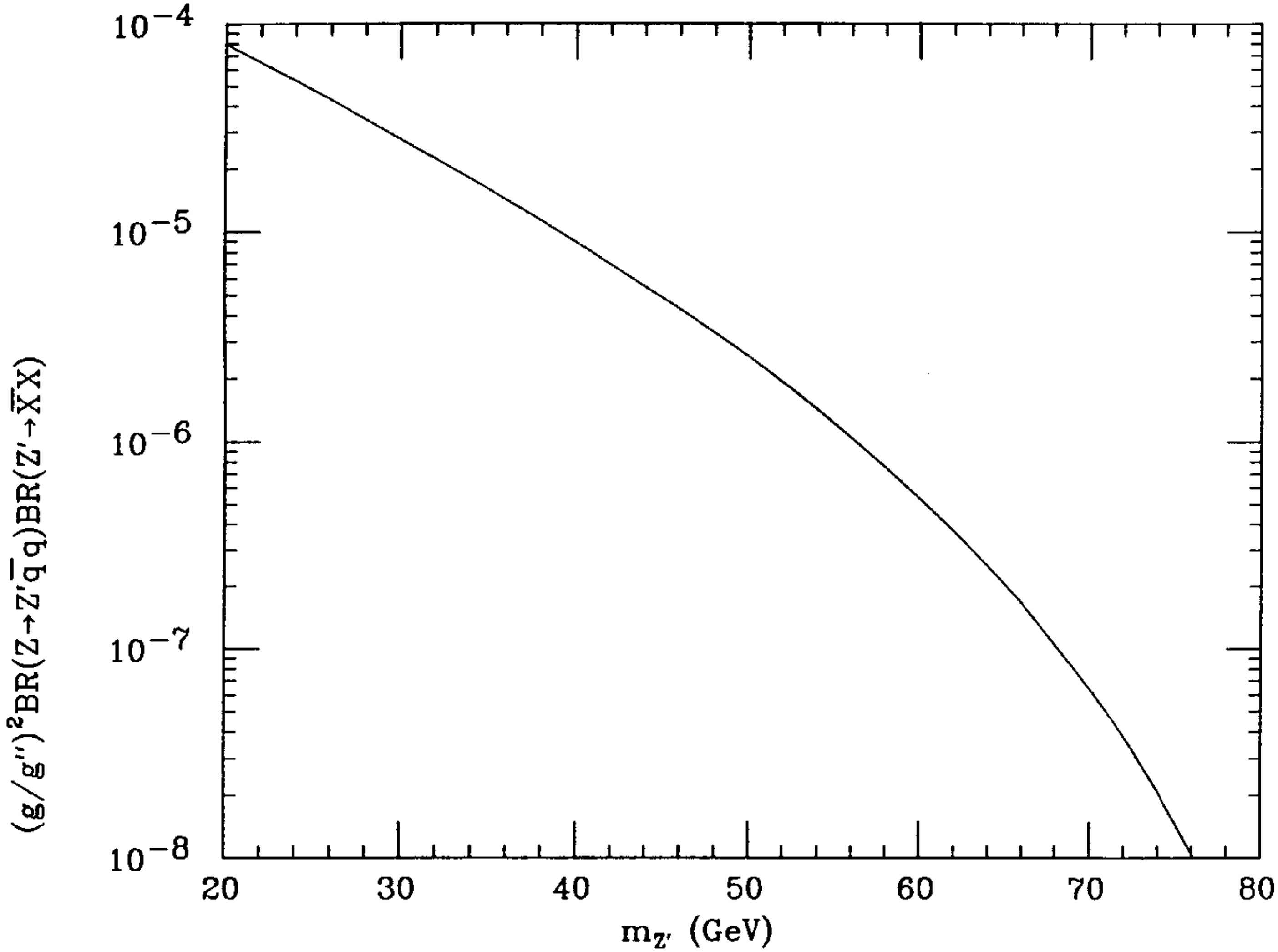
$(g/g'')^2 BR(Z \to Z'\bar{q}q) BR(Z' \to \bar{X}X)$
10^{-4}
10^{-5}
10^{-6}
10^{-7}
10^{-8}
20
30
40
50
60
70
80
$m_{Z'}$ (GeV)

Fig. 1

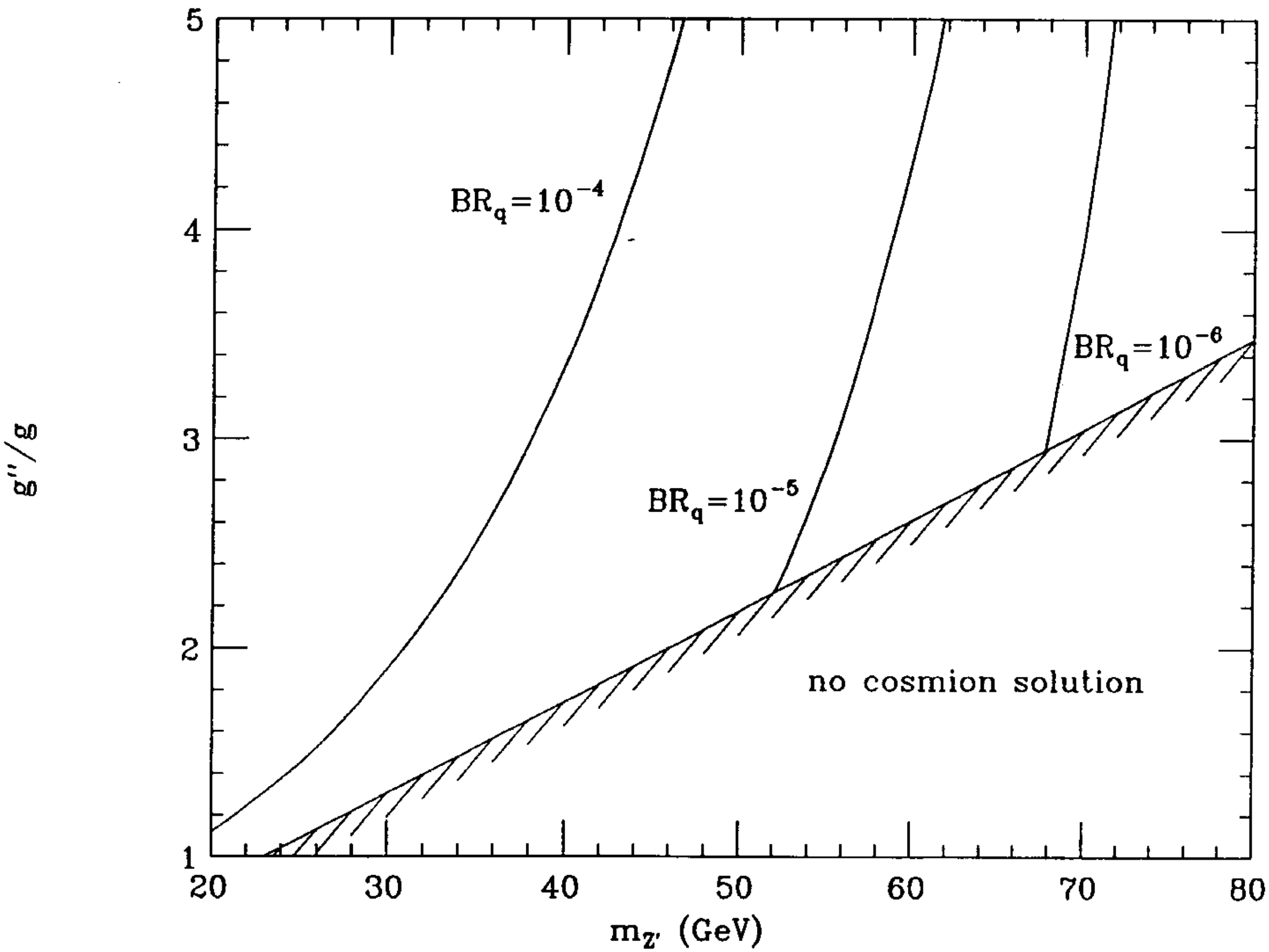
g''/g
5
4
3
2
1
20
30
40
50
60
70
80
m_Z' (GeV)
BR_q=10^-4
BR_q=10^-5
BR_q=10^-6
no cosmion solution

Fig. 2