PHYSICAL REVIEW D VOLUME 53, NUMBER 6 15 MARCH 1996

Light photinos as dark matter

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(Received 27 April 1995)

There are good reasons to consider models of low-energy supersymmetry with very light photinos and gluinos. In most of these models the lightest R-odd, color-singlet state containing a gluino, the R^0, has a mass in the 1–2 GeV range and the slightly lighter photino $\tilde{\gamma}$ would survive as the relic R-odd species. For the light photino masses considered here, previous calculations resulted in an unacceptable photino relic abundance. But we point out that processes other than photino self-annihilation determine the relic abundance when the photino and R^0 are close in mass. Including $R^0 \leftrightarrow \tilde{\gamma}$ processes, we find that the photino relic abundance is most sensitive to the R^0-to-$\tilde{\gamma}$ mass ratio, and within model uncertainties, a critical density in photinos may be obtained for an R^0-to-$\tilde{\gamma}$ mass ratio in the range 1.2 to 2.2. We propose photinos in the mass range of 500 MeV to 1.6 GeV as a dark matter candidate, and discuss a strategy to test the hypothesis.

PACS number(s): 98.80.Cq, 11.30.Pb, 14.80.Ly

I. INTRODUCTION

In this paper we study the early-Universe evolution and freeze-out of light, long-lived or stable, R-odd states, the photinos $\tilde{\gamma}$ and the gluino $\tilde{g}$.[1] In the type of models we consider, the photino should be the relic R-odd particle, even though it may be more massive than the gluino. This is because below the confinement transition the gluino is bound with a gluon into a color-singlet hadron, the R^0, whose mass (which is in the 1 to 2 GeV range when the gluino is very light [1,2]) is greater than that of the photino. Including previously neglected reactions associated with the gluino (more precisely, associated with the R^0), we find that light photinos may be cosmologically acceptable; indeed they are an attractive dark-matter candidate.

In the minimal supersymmetric (SUSY) model, the mass matrix of the charged and neutral SUSY fermions (gauginos and Higgsinos) are determined by Lagrangian terms involving the Higgs chiral superfields $\tilde{H}_1$ and $\tilde{H}_2$ and the SU(2) and U(1) gauge superfields $\tilde{W}^a$ and $\tilde{B}$ plus soft supersymmetry-breaking terms. This leads to a neutralino mass matrix in the basis $(\tilde{B},\tilde{W}^3,\tilde{H}_1^0,\tilde{H}_2^0)$ of the form

$$\begin{pmatrix} M_1 & 0 & -m_Z\cos\beta\,\sin\theta_W & m_Z\sin\beta\,\sin\theta_W \\ 0 & M_2 & m_Z\cos\beta\,\cos\theta_W & -m_Z\sin\beta\,\cos\theta_W \\ -m_Z\cos\beta\,\sin\theta_w & m_Z\cos\beta\,\cos\theta_W & 0 & -\mu \\ m_Z\sin\beta\,\sin\theta_W & -m_Z\sin\beta\,\cos\theta_W & -\mu & 0 \end{pmatrix}. \tag{1}$$

Here m_Z is the mass of the Z boson, θ_W is the Weinberg angle, μ is the coefficient of a supersymmetric mixing term between Higgs superfields, and $\tan\beta$ is the ratio of the vacuum expectation values of the two Higgs fields responsible for electroweak symmetry breaking. The SUSY-breaking masses M_1 and M_2 are commonly assumed to be of order m_Z or larger, and if the SUSY model is embedded in a grand unified theory, then $3M_1/M_2=5\alpha_1/\alpha_2$.

The terms in the Lagrangian proportional to M_1 and M_2 arise from dimension-3 SUSY-breaking operators. However such SUSY-breaking terms are not without problems. It appears difficult to break SUSY dynamically in a way that produces dimension-3 terms while avoiding problems associated with the addition of gauge-singlet superfields [3]. In models where SUSY is broken dynamically or spontaneously

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[1] R parity is a multiplicative quantum number, exactly conserved in most SUSY models, under which ordinary particles have $R=+1$ while new "superpartners" have $R=-1$. Throughout this paper we will assume that R parity is exact so the lightest R-odd particle is stable.

0556-2821/96/53(6)/2990(12)/$10.00

in the hidden sector and there are no gauge singlets, all dimension-3 SUSY-breaking operators in the effective low-energy theory are suppressed compared to SUSY breaking scalar masses by a factor of $\langle\Phi\rangle/m_{\rm Pl}$, where $\langle\Phi\rangle$ is the vacuum expectation value of some hidden-sector field. Thus, dimension-3 terms do not contribute to the low-energy effective Lagrangian. This would imply that at the tree level the gluino is massless, and the neutralino mass matrix is given by Eq. (1) with vanishing (00) and (11) entries. However, nonzero contributions to the gluino mass and the neutralino mass matrix come from two sources: radiative corrections such as the top-quark–top-squark loops for the gluino and neutralinos, and "electroweak" loops involving Higgsinos and/or W-inos and B-inos for the neutralinos (but not for the gluino).

The generation of radiative gaugino masses in the absence of dimension-3 SUSY breaking was studied by Farrar and Masiero [4].[2] From Figs. 4 and 5 of that paper one sees (taking $\mu \gtrsim 40$ GeV) that as the typical SUSY-breaking scalar mass M_0 varies between 100 and 400 GeV, the gluino mass ranges from about 700 to about 100 MeV,[3] while the photino[4] mass ranges from around 400 to 900 MeV. This estimate for the photino mass should be considered as merely indicative of its possible value, since an approximation for the electroweak loop used in Ref. [4] is strictly valid only when μ and M_0 are much larger than m_W. The other neutralinos are much heavier, and the production rates of the photino and the next-lightest neutralino in Z^0 decay are consistent with LEP bounds [4].

Using the results of Ref. [4], but additionally restricting parameters so that the correct electroweak symmetry breaking is obtained, Farrar [2] found $M_0 \sim 150$ GeV and estimated the R^0 lifetime. This allowed completion of the study of the main phenomenological features of this scenario, which was begun in Ref. [1]. The conclusion is that light gluinos and photinos are quite consistent with present experiments, and result in a number of striking predictions [2]. However, models with light gauginos have been widely thought to be disallowed because it has been believed that the relic density of the lightest neutralino, usually referred to as the lightest supersymmetric partner (LSP),[5] exceeds cosmological bounds unless R-parity is violated [8,9].

In this paper we point out that previous considerations of the relic abundance have neglected the rather important interplay between the photino and the gluino which can determine the final neutralino abundance if the photino and gluino are both light, as they must be in models without dimension-3 explicit SUSY-breaking terms. We find that when gluino–photino interactions are included, rather than being a cosmological embarrassment, these very light photinos are an excellent dark matter candidate. In this paper we discuss the decoupling and relic abundance of light photinos, and the sensitivity of the result upon the parameters of the SUSY models.

For the light masses studied here, freeze-out occurs well after the confinement transition so the physical states must be color singlets. Since $\tilde{g}$ is not a color singlet, below the confinement transition the relevant state to consider is the lightest color-singlet state containing a gluino, which is believed to be a gluon–gluino bound state known as the R^0. The other light R-odd states are more massive than these, and decay to the two light ones with lifetimes much faster than the expansion rate at freeze-out. The only other possible state of interest is the S^0, which is the lightest R-odd baryon, consisting of the color-singlet, flavor-singlet state $uds\tilde{g}$ [1,2,10].

Since the $\tilde{\gamma}$ is the lightest *color-singlet* R-odd state, it is stable. The R^0 decays to a final state consisting of a photino and typically one meson: $R^0 \to \tilde{\gamma}\pi, \tilde{\gamma}\eta$, etc. The lifetime is very uncertain, but probably lies in the range 10^{-4} to 10^{-10} s, or even longer [2].

While the predictions of Ref. [2] for R hadron and photino masses motivated the present work, the analysis we give here is general and applies for other mass ranges as well. We will see that the result is very sensitive to the *ratio* of R^0 and photino masses, but rather insensitive to other parameters. However, for concreteness the reader may wish to keep in mind the mass ranges predicted in [2,4]:

$$\begin{aligned} \tilde{g}(\text{gluino}):&\quad m_{\tilde{g}} = 100-600\ \text{MeV},\\ \tilde{\gamma}(\text{photino}):&\quad m = 100-1400\ \text{MeV},\\ R^0(\tilde{g}g):&\quad M = 1-2\ \text{GeV},\\ S^0(uds\tilde{g}):&\quad M_{S^0} = 1.5\text{–}2\ \text{GeV}. \end{aligned} \tag{2}$$

The reaction rates that determine freeze-out will depend upon the $\tilde{\gamma}$ and R^0 masses, the cross sections involving the $\tilde{\gamma}$ and R^0, and possibly the decay width of the R^0 as well. In turn the cross sections and decay width also depend on the masses of the $\tilde{\gamma}$, $\tilde{g}$, and R^0, as well as the masses of the squarks and sleptons. We will denote the squark or slepton masses by a common mass scale $M_{\tilde{S}}$ (expected to be of order 100 GeV). Even if the masses were known and the short-distance physics specified, calculation of the width and some of the cross sections would be no easy task, because one is dealing with light hadrons. Fortunately, our conclusions are reasonably insensitive to individual masses, lifetimes, and cross sections, but depend crucially upon the R^0-to-$\tilde{\gamma}$ mass ratio. When we do need an explicit value of the photino mass

[2] See also [5] for general formulas. Earlier studies [6,7] of radiative corrections when tree-level gaugino masses are absent included another dimension-3 operator, the "A term." They also assumed model-dependent relations between parameters.

[3] Gluino masses quoted in [4] are values at the electroweak scale; at the hadronic scale the mass is larger due to renormalization group running. Smaller μ and larger M_0 lead to smaller values of the gluino mass. However 100 MeV is still the operative order of magnitude of the lower limit on the gluino mass at the hadronic scale because a sufficiently light gluino at the hadronic scale leads to an unacceptably light pseudoscalar meson [1].

[4] Upon diagonalization of the mass matrix, the physical neutralino states are a linear combination of $\tilde{B}^0$, $\tilde{W}^3$, $\tilde{H}_1^0$, and $\tilde{H}_2^0$. When the gaugino submatrix elements are small, the lightest neutralino is a linear combination of $\tilde{W}^3$ and $\tilde{B}^0$ that is almost identical to the SU(2)×U(1) composition of the photon, and thus is correctly called "photino."

[5] In this scenario, LSP is an ambiguous term: the gluino is lighter than the photino, although the photino is lighter than the R^0. A more relevant term would be LROCS—*lightest R-odd color singlet.*

m or the masses of squarks and sleptons, we will parametrize them by the dimensionless ratios

$$\mu_8 \equiv \frac{m}{800\ \text{MeV}}, \quad \mu_S \equiv \frac{M_{\tilde{S}}}{100\ \text{GeV}}. \tag{3}$$

Although there are several undetermined parameters in our calculation, as mentioned above, the most important parameter will be the ratio of the R^0 mass to the $\tilde{\gamma}$ mass:

$$r \equiv \frac{M}{m}. \tag{4}$$

This is by far the most crucial parameter, with the relic abundance having an exponential dependence upon r. We find that limits to the magnitude of the contribution to the present mass density from relic photinos requires[6] $r \lesssim 2.2$, while r must be larger than about 1.2 if the photino relic density is to be significant. This narrow band of r encompasses the large uncertainties in lifetimes and cross sections. If the mass ratio is between about 1.6 and 2, then light-mass relic photinos dominate the Universe and provide the dark matter with $\Omega_{\tilde{\gamma}} \sim 1$.

In the concluding section we explore the proposal that light photinos are the dark matter, and discuss possibilities for testing the idea. We lay the groundwork for this suggestion in the next section as we develop a new scenario for decoupling and freeze-out for the photinos and gluinos. In Sec. III we consider the cross sections and lifetimes used in Sec. IV to calculate the reaction rates relevant for the determination of the freeze-out abundance of the photinos (and hence $\Omega_{\tilde{\gamma}} h^2$). In Sec. V we compare the reaction rates to the expansion rate and estimate when photinos decouple.

II. SCENARIO FOR PHOTINO OR GLUINO FREEZE-OUT

The standard procedure for the calculation of the present number density of a thermal relic of the early-Universe is to assume that the particle species was once in thermal equilibrium until at some point the rates for self-annihilation and pair-creation processes became much smaller than the expansion rate, and the particle species effectively froze out of equilibrium. After freeze-out, its number density decreased only because of the dilution due to the expansion of the Universe. (For a discussion, see Ref. [11].)

Since after freeze-out the number of particles in a *comoving* volume is constant, it is convenient to express the number density of the particle species in terms of the entropy density, since the entropy in a comoving volume is also constant for most of the history of the Universe. The number-density-to-entropy ratio is usually denoted by Y. If a species of mass m is in equilibrium and nonrelativistic, Y is simply given in terms of the mass-to-temperature ratio $x \equiv m/T$ as

$$Y_{\text{eq}}(x) = 0.145(g/g_*)x^{3/2}\exp(-x), \tag{5}$$

where g is the number of spin degrees of freedom, and g_* is the total number of relativistic degrees of freedom in the Universe at temperature $T = m/x$. Well after freeze-out $Y(x)$ is constant, and we will denote this asymptotic value of Y as Y_∞.

If self-annihilation determines the final abundance of a species, Y_∞ can be found by integrating the Boltzmann equation (an overdot denotes d/dt)

$$\dot{n} + 3Hn = -\langle |v| \sigma_A \rangle (n^2 - n_{\text{eq}}^2), \tag{6}$$

where n is the actual number density, n_{eq} is the equilibrium density, H is the expansion rate of the Universe, and $\langle |v| \sigma_A \rangle$ is the thermal average [12,13] of the annihilation rate.

There are no general closed-form solutions to the Boltzmann equation, but there are reliable, well tested approximations for the late-time solution, i.e., Y_∞. Then with knowledge of Y_∞, the contribution to Ωh^2 from the species can easily be found. Let us specialize to the survival of photinos assuming self-annihilation determines freeze out.

Calculation of the relic abundance involves first calculating the value of x, known as x_f, where the abundance starts to depart from the equilibrium abundance. Using standard approximate solutions to the Boltzmann equation [11] gives[7]

$$x_f = \ln(0.0481 m_{\text{Pl}} m \sigma_0) - 1.5\ \ln[\ln(0.0481 m_{\text{Pl}} m \sigma_0)], \tag{7}$$

where we have used $g = 2$ and $g_* = 10$, and parametrized the nonrelativistic annihilation cross section as $\langle |v| \sigma_A \rangle = \sigma_0 x^{-1}$. In anticipation of the results of the next section, we use $\sigma_0 = 2 \times 10^{-11} \mu_8^2 \mu_S^{-4}$ mb, and we find $x_f \simeq 12.3 + \ln(\mu_8^3/\mu_S^4)$. The value of x_f determines Y_∞:

$$Y_\infty = \frac{2.4 x_f^2}{m_{\text{Pl}} m \sigma_0} \simeq 7.4 \times 10^{-7} \mu_8^{-3} \mu_S^4. \tag{8}$$

Once Y_∞ is known, the present photino energy density can be easily calculated: $\rho_{\tilde{\gamma}} = m n_{\tilde{\gamma}} = 0.8 \mu_8$ GeV$\times Y_\infty$ 2970 cm^{-3}. When this result is divided by the critical density, $\rho_C = 1.054 h^2 \times 10^{-5}$ GeV cm^{-3}, the fraction of the critical density contributed by the photino is $\Omega_{\tilde{\gamma}} h^2 = 2.25 \times 10^8$ $\mu_8 Y_\infty$. For Y_∞ in Eq. (8), $\Omega_{\tilde{\gamma}} h^2 = 167 \mu_8^{-2} \mu_S^4$.

The age of the Universe restricts $\Omega_{\tilde{\gamma}} h^2$ to be less than one, so for $\mu_S = 1$, the photino must be more massive than about 10 GeV if its relic abundance is determined by self-annihilation.

But in this paper we point out that for models in which both the photino and the gluino are light, freeze-out is not determined by photino self-annihilation, but by $\tilde{\gamma}$–R^0 interconversion. The basic point is that since the R^0 has strong interactions, it will stay in equilibrium longer than the photino, even though it is more massive. As long as $\tilde{\gamma} \leftrightarrow R^0$ interconversion occurs at a rate larger than H, then through its interactions with the R^0 the photino will be able to maintain its equilibrium abundance even after self-annihilation has

[6]Or else R parity must be violated so the photinos decay.

[7]Freeze-out aficionados will notice that we use the formulas appropriate for p-wave annihilation because Fermi statistics requires the initial identical Majorana fermions to be in an $L = 1$ state [8].

frozen out.[8]

Griest and Seckel discussed the possibility that the relic abundance of the lightest species is determined by its interactions with another species [14]. They concluded that the mass splitting between the relic and the heavier particle must be less than 10% for the effect to be appreciable. We find that R^0-$\tilde{\gamma}$ interconversion determines the $\tilde{\gamma}$ relic abundance even though the R^0 may be twice as massive as the $\tilde{\gamma}$. The difference arises because Griest and Seckel assumed that all cross sections were roughly the same order of magnitude. But in our case the R^0 annihilation is about 10^{12} times larger than other relevant cross sections.

Before we demonstrate that this scenario naturally occurs for the types of photino and R^0 masses expected, we must determine the cross sections and decay width of the reactions involving the photino and the gluino.

III. CROSS SECTIONS AND DECAY WIDTH

In this section we characterize the cross sections and decay width required for the determination of the relic photino abundance, and also discuss the uncertainties. We should emphasize that all cross sections are calculated in the nonrelativistic (NR) limit, and by $\langle\cdots\rangle$ we imply that the quantity is to be evaluated as a thermal average [12,13]. In the NR limit a temperature dependence to the cross sections enters if the annihilation proceeds through a p wave, as required if the initial state consists of identical fermions [8]. For p-wave annihilation, at low energy the cross section is proportional to v^2, where v is the relative velocity of the initial particles. The thermal average reduces to replacing v^2 by $6T/m$, where m is the mass of the particle in the initial state.

We now consider in turn the cross sections and width for the individual reactions discussed in the previous section.

A. Self-annihilations and coannihilation

The first type of reactions we will consider are those which change the number of R-odd particles.

$R^0R^0\to X$. We will refer to this process as R^0 self-annihilation. At the constituent level the relevant reactions are $\tilde{g}+\tilde{g}\to g+g$ and $\tilde{g}+\tilde{g}\to q+\bar{q}$, which are unsuppressed by any powers of $M_{\tilde{S}}$, and should be typical of strong interaction cross sections. In the NR limit, we expect the R^0R^0 annihilation cross section to be comparable to the $\bar{p}p$ cross section, but with an extra factor of v^2, accounting for the fact that there are identical fermions in the initial state, so annihilation must proceed through a p wave.[9] There is some energy dependence to the $\bar{p}p$ cross section, but it is sufficient to consider $\langle|v|\sigma_{R^0R^0}\rangle$ to be a constant, approximately given by

$$\langle|v|\sigma_{R^0R^0}\rangle\simeq 100v^2\ \ \text{mb}=600x^{-1}r^{-1}\ \ \text{mb}, \tag{9}$$

where we have used for the relative velocity $v^2=6T/M=6/(rx)$, with $x\equiv m/T$.

We should note that the thermal average of the cross section might be even larger if there are resonances near threshold. In any case, this cross section should be much larger than any cross section involving the photino, and will ensure that the R^0 remains in equilibrium longer than the $\tilde{\gamma}$, greatly simplifying our considerations.

$\tilde{\gamma}\tilde{\gamma}\to X$. In photino self-annihilation at low energies the final state X is a lepton-antilepton pair, or a quark-antiquark pair which appears as light mesons. The process involves the t-channel exchange of a virtual squark or slepton between the photinos, producing the final-state fermion-antifermion pair. In the low-energy limit the mass $M_{\tilde{S}}$ of the squark or slepton is much greater than $\sqrt{s}$, and the photino-photino-fermion-antifermion operator appears in the low-energy theory with a coefficient proportional to $e_i^2/M_{\tilde{S}}^2$, with e_i the charge of the final-state fermion.[10] Also, as there are two identical fermions in the initial state, the annihilation proceeds as a p wave, which introduces a factor of v^2 in the low-energy cross section [8]. The resultant low-energy photino self-annihilation cross section is [8,9,15,16]

$$\begin{aligned}\langle|v|\sigma_{\tilde{\gamma}\tilde{\gamma}}\rangle&=8\pi\alpha_{\text{em}}^2\sum_i\ \ g_i^4\,\frac{m^2}{M_{\tilde{S}}^4}\,\frac{v^2}{3}\\&\simeq 2.0\times 10^{-11}x^{-1}[\mu_8^2\mu_S^{-4}]\ \ \text{mb},\end{aligned} \tag{10}$$

where we have used for the relative velocity $v^2=6/x$ with $x\equiv m/T$, and q_i is the magnitude of the charge of a final-state fermion in units of the electron charge. For the light photinos we consider, summing over e, μ, and three colors of u, d, and s quarks leads to $\Sigma_i q_i^4=8/3$.

$\tilde{\gamma}R^0\to X$. This is an example of a phenomenon known as coannihilation [14,17] whereby the particle of interest (in our case the photino) disappears by annihilating with another particle (here, the R^0). Of course coannihilation also leads to a net decrease in R odd particles.

In all processes involving the photino–R^0 interaction, the leading tree-level short-distance operator containing $\tilde{g}$ and $\tilde{\gamma}$ is $\lambda_{\tilde{g}}^{\dagger}\lambda_{\tilde{\gamma}}q_i^{\dagger}q_i+\text{H.c.}$, with coefficient $eq_ig_S/M_{\tilde{S}}^2$. For three

[8] Actually, interconversion can also play an important role in determining the relic abundance of heavier photinos. When the photino is more massive and freeze-out occurs above the confinement phase transition, the analysis is similar to the one here; in fact it is simpler because perturbation theory can be used to compute the relevant rates involving gluinos and photinos. Since the qualitative relation between interconversion and self-annihilation rates is independent of whether the gluino is free or confined in an R^0, one can get a crude idea of the required gluino-photino mass ratio, r, just by using the analysis in this paper and scaling the results to the value of μ_8 of interest. We concentrate on the light gaugino scenario because it is attractive in its own right, and also because it *naturally* produces r to a good approximation [2]. In a conventional SUSY-breaking scheme fine-tuning is generally necessary to give r the right value for the interconversion mechanism to play an important role.

[9] In general the result is not so simple. For instance, in addition to the term proportional to v^2, the cross section also involves a term proportional to the square of the masses of the initial and final particles.

[10] The electric charge e and the strong charge g_S are to be evaluated at a scale of order $M_{\tilde{S}}$, so in numerical estimates we use $\alpha_{\text{em}}=1/128$ and $\alpha_S=0.117$.

light quarks, $\Sigma_i q_i^2=2$. Thus we can estimate the cross section for $\tilde{\gamma}R^0 \to X$ in terms of the $\tilde{\gamma}$ self-annihilation cross section:

$$\langle |v| \sigma_{\tilde{\gamma}R^0} \rangle \simeq \frac{\alpha_S}{\alpha_{\rm em}} \frac{4}{3} \frac{2}{8/3} \frac{M}{m} \frac{3}{v^2} \langle |v| \sigma_{\tilde{\gamma}\tilde{\gamma}} \rangle, \tag{11}$$

where the ratio of α's arises because the short-distance operator for coannihilation is proportional to $e_i^2 g_S^2$ rather than e_i^4, the second factor is the color factor coming from the gluino coupling, and the third factor comes from the ratio of $\Sigma_i q_i^2 / \Sigma_i q_i^4$ for the participating fermions. We have replaced m^2 appearing in Eq. (10) by mM, although the actual dependence on m and M may be more complicated. Finally, the annihilation is s wave so there is no $v^2/3$ suppression as in photino self-annihilation.

Although the short-distance physics is perturbative, the initial gluino appears in a light hadron, and there are complications in the momentum fraction of the R^0 carried by the gluino and other nonperturbative effects. For our purposes it will be sufficient to account for the uncertainty by including in the cross section and unknown coefficient A, leading to a final expression

$$\langle |v| \sigma_{\tilde{\gamma}R^0} \rangle \simeq 1.5 \times 10^{-10} r [\mu_8^2 \mu_S^{-4} A] \ \text{mb}. \tag{12}$$

It is reassuring that if one estimates $\langle |v| \sigma_{\tilde{\gamma}R^0} \rangle$ starting from $\langle |v| \sigma_{R^0 R^0} \rangle$ a similar result is obtained. We find that coannihilation will not be important unless A is larger than 10^2 or so, which we believe is unlikely.

B. $\tilde{\gamma}$–R^0 interconversion

In what we call interconversion processes, there is an R-odd particle in the initial as well as the final state. Although the reactions do not of themselves change the number of R-odd particles, they keep the photinos in equilibrium with the R^0's, which in turn are kept in equilibrium through their self-annihilations.

$R^0 \to \tilde{\gamma}\pi$. R^0 decay can occur via, e.g., the gluino inside the initial R^0 turning into an antiquark and a virtual squark, followed by squark decay into a photino and a quark. In the low-energy limit the quark-antiquark–gluino-photino vertex can be described by the same type of four-Fermi interaction as in coannihilation. One expects on dimensional grounds a decay width $\Gamma_0 \propto \alpha_{\rm em} \alpha_S M^5 / M_{\tilde{S}}^4$. The lifetime of a free gluino to decay to a photino and massless quark-antiquark pair was computed in Ref. [18]. However, this does not provide a very useful estimate when the gluino mass is less than the photino mass.

The lifetime for R^0 decay was studied in Ref. [2]. In an attempt to account for the effects of gluino-gluon interactions in the R^0, necessary for even crude estimate of the R^0 lifetime, the following picture was developed, based on the approach of Altarelli *et al.* [19]: The R^0 is viewed as a state with a massless gluon carrying momentum fraction x, and a gluino carrying momentum fraction $(1-x)$,[11] having therefore an effective mass $M\sqrt{1-x}$. The gluon structure function $F(x)$ gives the probability in an interval x to $x+dx$ of finding a gluon, and the corresponding effective mass for the gluino. One then obtains the R^0 decay width (neglecting the mass of final-state hadrons)

$$\Gamma(M,r) = \Gamma_0(M,0) \int_0^{1-r^{-2}} dx (1-x)^{5/2} F(x) f(1/r\sqrt{1-x}), \tag{13}$$

where $\Gamma_0(M,0)$ is the rate for a gluino of mass M to decay to a massless photino, and $f(y) = [(1-y^2)(1+2y-7y^2+20y^3-7y^4+2y^5+y^6)+24y^3(1-y+y^2)\log y]$ contains the phase-space suppression which is important when the photino becomes massive in comparison to the gluino. Modeling $K^\pm$ decay in a similar manner underestimates the lifetime by a factor of 2 to 4. This is in surprisingly good agreement; however, caution should be exercised when extending the model to R^0 decay, because kaon decay is much less sensitive to the phase-space suppression from the final-state masses than the present case, since the range of interest will turn out to be $r \sim 1.2$–2.2. For r in this range, taking $F(x) \sim 6x(1-x)$ following the discussion in Ref. [2] leads to an approximate behavior

$$\Gamma_{R^0 \to \tilde{\gamma}\pi} = 2.0 \times 10^{-14} \mathscr{F}(r) \ \text{GeV} [\mu_8^5 \mu_S^{-4} B], \tag{14}$$

where $\mathscr{F}(r) = r^5 (1-r^{-1})^6$, and the factor B reflects the overall uncertainty. We believe a reasonable range for B is $1/30 \lesssim B \lesssim 3$. Lattice QCD calculation of the relevant hadronic matrix elements would allow a more reliable determination of B.

$R^0 \pi \leftrightarrow \tilde{\gamma}\pi$. We will refer to these processes as photino-R_0 conversion, since an initial R^0 (or $\tilde{\gamma}$) is converted to a final $\tilde{\gamma}$ (or R^0). The short-distance subprocess in this reaction is $q+\tilde{g} \to q+\tilde{\gamma}$, again described by the same low-energy effective operator as in coannihilation and R^0 decay. At the hadronic level the matrix element for $R^0 \pi \to \tilde{\gamma} X$ is the same as for $R^0 \tilde{\gamma} \to \pi X$ for any X, evaluated in different physical regions. Thus the difference between the various cross sections is just due to the difference in fluxes and final-state phase-space integrations, and variations of the matrix element with kinematic variables. Given the crude nature of the analysis here, and the great uncertainty in the overall magnitude of the cross sections, incorporating the constraints of crossing symmetry are not justified at present. We will therefore use the same form as for Eq. (12), letting C parametrize the hadronic uncertainty in this case:

$$\langle |v| \sigma_{R^0 \pi} \rangle \simeq 1.5 \times 10^{-10} r [\mu_8^2 \mu_S^{-4} C] \ \text{mb}. \tag{15}$$

We can point to one specific hadronic effect which is not explicitly included in (15), but which is potentially important. It is likely that near threshold there is a resonance (the R_π) which would increase the cross section by a factor of $4M_R^2/\Gamma_R^2$, where M_R is the mass and Γ_R the width of the resonance. This complicates matters because neither the resonance's width nor its distance above threshold is known. If a resonance is important, it would also be necessary to

[11] Of course there should be no confusion with the fact that in the discussion of the R^0 lifetime we use x to denote the gluon momentum fraction whereas throughout the rest of the paper x denotes m/T.

TABLE I. Cross sections and the decay width used in the calculation of the relic photino abundance. The dimensionless parameters μ_8 and μ_S were defined in Eq. (3), and $\mathscr{F}(r)$ was discussed below Eq. (14). The coefficients A, B, and C reflect uncertainties involving the calculation of hadronic matrix elements.

Process		Cross section or width
R^0 self-annihilation:	$\langle \lvert v\rvert \sigma_{R^0R^0}\rangle$	$600x^{-1}r^{-1}$ mb
$\tilde{\gamma}$ self annihilation:	$\langle \lvert v\rvert \sigma_{\tilde{\gamma}\tilde{\gamma}}\rangle$	$2.0\times10^{-11}x^{-1}[\mu_8^2\mu_S^{-4}]$ mb
Coannihilation:	$\langle \lvert v\rvert \sigma_{\tilde{\gamma}R^0}\rangle$	$1.5\times10^{-10}r[\mu_8^2\mu_S^{-4}A]$ mb
R^0 decay:	$\Gamma_{R^0\rightarrow\tilde{\gamma}\pi}$	$2.0\times10^{-14}\mathscr{F}(r)[\mu_8^5\mu_S^{-4}B]$ GeV
$\tilde{\gamma}$–R^0 conversion:	$\langle \lvert v\rvert \sigma_{R^0\pi}\rangle$	$1.5\times10^{-10}r[\mu_8^2\mu_S^{-4}C]$ mb

perform the thermal average over the resonance in a more careful manner [13]. In order to take into account the possibility of such a resonance, we allow C to vary in the range $1\lesssim C\lesssim10^3$. We will use detailed balance arguments which allow us to avoid using the inverse reaction, $\tilde{\gamma}\pi\rightarrow R^0\pi$.

This completes the discussion of the lifetimes, cross sections, and their uncertainties. The results are summarized in Table I.

IV. EARLY-UNIVERSE REACTION RATES

To obtain an estimate of when the rates will drop below the expansion rate, we will assume all particles are in local thermodynamic equilibrium (LTE). In LTE a particle of mass m in the NR limit has a number density

$$n=\frac{g}{(2\pi)^{3/2}}\,(mT)^{3/2}\exp(-m/T)$$

$$=\frac{g}{(2\pi)^{3/2}}\,(T/m)^{3/2}m^3\exp(-m/T). \qquad (16)$$

Here g counts the number of spin degrees of freedom, and will be 2 for the R^0 and the $\tilde{\gamma}$.

H (the expansion rate). Of course all rates are to be compared with the expansion rate. In the radiation-dominated Universe with $g_*\sim10$ degrees of freedom

$$H=1.66g_*^{1/2}T^2/m_{\rm Pl}=2.8\times10^{-19}x^{-2}[\mu_8^2]\ \ \text{GeV}. \qquad (17)$$

$\tilde{\gamma}\tilde{\gamma}\rightarrow X$ *(photino self-annihilation).* In the Boltzmann equation for the evolution of the $\tilde{\gamma}$ number density there are terms accounting for photino self-annihilation and photino pair production from light particles in the plasma. Assuming the light annihilation products are in LTE, the terms are of the form

$$\dot{n}_{\tilde{\gamma}}+3Hn_{\tilde{\gamma}}\supset-\langle|v|\sigma_{\tilde{\gamma}\tilde{\gamma}}\rangle[(n_{\tilde{\gamma}})^2-(n_{\tilde{\gamma}}^{\rm eq})^2]. \qquad (18)$$

If we assume that the photino is in equilibrium, the self-annihilation and pair production terms are equal, and we may express the individual terms in the form

$$\dot{n}_{\tilde{\gamma}}\supset-3Hn_{\tilde{\gamma}}\mp[\langle|v|\sigma_{\tilde{\gamma}\tilde{\gamma}}\rangle n_{\tilde{\gamma}}^{\rm eq}]n_{\tilde{\gamma}}^{\rm eq}, \qquad (19)$$

where the upper sign is for self-annihilation and the lower sign is for pair production.

It is obvious that $[n_{\tilde{\gamma}}^{\rm eq}\langle|v|\sigma_{\tilde{\gamma}\tilde{\gamma}}\rangle]$ plays the role of a ''rate'' to be compared to H. If this rate is much greater than H, the self-annihilation or pair-production processes will ensure the photino is in equilibrium, while if the rate is much smaller than H, self-annihilation or pair production cannot enforce equilibrium.

Therefore we define an equilibrium photino annihilation rate by $\Gamma(\tilde{\gamma}\tilde{\gamma}\rightarrow X)=n_{\tilde{\gamma}}^{\rm eq}\langle|v|\sigma_{\tilde{\gamma}\tilde{\gamma}}\rangle$. Using Eq. (16) for the equilibrium abundance and the annihilation cross section discussed in the previous section, we find

$$\Gamma(\tilde{\gamma}\tilde{\gamma}\rightarrow X)=\frac{2}{(2\pi)^{3/2}}\left(\frac{T}{m}\right)^{3/2}m^3$$

$$\times\exp(-m/T)\ \frac{2.0\times10^{-11}\ \text{mb}}{0.39\ \text{mb GeV}^2}\ x^{-1}[\mu_8^2\mu_S^{-4}]$$

$$=3.3\times10^{-12}x^{-5/2}\exp(-x)[\mu_8^5\mu_S^{-4}]\ \ \text{GeV}. \qquad (20)$$

$R^0R^0\rightarrow X$ (R^0 *self-annihilation).* Determination of the equilibrium rate for R^0 self-annihilation proceeds in a similar manner, yielding $\Gamma(R^0R^0\rightarrow X)=n_{R^0}^{\rm eq}\langle|v|\sigma_{R^0R^0}\rangle$:

$$\Gamma(R^0R^0\rightarrow X)=\frac{2}{(2\pi)^{3/2}}\left(\frac{T}{M}\right)^{3/2}M^3$$

$$\times\exp(-M/T)\ \frac{240x^{-1}r^{-1}\ \text{mb}}{0.39\ \text{mb GeV}^2}$$

$$=99r^{1/2}x^{-5/2}\exp(-rx)[\mu_8^3]\ \ \text{GeV}. \qquad (21)$$

$\tilde{\gamma}R^0\rightarrow X$ ($\tilde{\gamma}$ *coannihilation).* In the Boltzmann equation for the evolution of the $\tilde{\gamma}$ density will appear a term $-n_{R^0}n_{\tilde{\gamma}}\langle|v|\sigma_{\tilde{\gamma}R^0}\rangle$. Therefore the equilibrium coannihilation rate for the decrease of the $\tilde{\gamma}$ density is

$$\Gamma(\tilde{\gamma}R^0\rightarrow X)=n_{R^0}^{\rm eq}\langle|v|\sigma_{\tilde{\gamma}R^0}\rangle$$

$$=2.5\times10^{-11}r^{5/2}x^{-3/2}\exp(-rx)[\mu_8^5\mu_S^{-4}], \qquad (22)$$

where we have again assumed the particles in the process are in equilibrium.

$\tilde{\gamma}\pi\rightarrow R^0$ *(inverse decay).* If the R^0 decay products (in this case $\tilde{\gamma}$ and π) are in equilibrium, then the Boltzmann equation for the evolution of R^0 contains a term

$$\dot{n}_{R^0}+3Hn_{R^0}\supset-\Gamma_{R^0\rightarrow\tilde{\gamma}\pi}(n_{R^0}-n_{R^0}^{\rm eq}). \qquad (23)$$

The first term on the right-hand side (RHS) represents decay, while the second term represents ''inverse decay.'' Since in-

TABLE II. The ratio of the equilibrium rates to the expansion rate for the indicated reactions. Shown in $[\cdots]$ is the scaling of the rates with unknown parameters characterizing the cross sections and decay width.

Process		Γ/H	
$\tilde{\gamma}$ self-annihilation	$\tilde{\gamma}\tilde{\gamma}\rightarrow X$	$1.2\times10^{7}x^{-1/2}\exp(-x)$	$[\mu_8^3\mu_S^{-4}]$
R^0 self-annihilation	$R^0R^0\rightarrow X$	$3.5\times10^{20}x^{-1/2}r^{1/2}\exp(-rx)$	$[\mu_8]$
Coannihilation	$\tilde{\gamma}R^0\rightarrow X$	$8.9\times10^{7}x^{1/2}r^{5/2}\exp(-rx)$	$[\mu_8^3\mu_S^{-4}A]$
Inverse decay	$\tilde{\gamma}\pi\rightarrow R^0$	$7.1\times10^{4}x^{2}r^{3/2}\mathscr{F}(r)\exp[-(r-1)x]$	$[\mu_8^3\mu_S^{-4}B]$
$\tilde{\gamma}$–R^0 conversion	$\tilde{\gamma}\pi\rightarrow R^0\pi$	$9.6\times10^{6}x^{1/2}r^{5/2}\exp[-(r-1)x]\exp(-0.175\mu_8^{-1}x)$	$[\mu_8^{3/2}\mu_S^{-4}C]$

verse decay turns a $\tilde{\gamma}$ into an R^0, there will be an "inverse decay" term in the equation for the evolution of the $\tilde{\gamma}$ number density:

$$\dot{n}_{\tilde{\gamma}}+3Hn_{\tilde{\gamma}}\supset-\Gamma_{R^0\rightarrow\tilde{\gamma}\pi}\,n_{R^0}^{\rm eq}. \tag{24}$$

The RHS can be written as $n_{\tilde{\gamma}}^{\rm eq}(n_{R^0}^{\rm eq}/n_{\tilde{\gamma}}^{\rm eq})\Gamma_{R^0\rightarrow\tilde{\gamma}\pi}$. Therefore the inverse decay rate in the evolution of the photino number density contributes a term

$$\begin{aligned}\Gamma(\tilde{\gamma}\pi\rightarrow R^0)&=\Gamma_{R^0\rightarrow\tilde{\gamma}\pi}(n_{R^0}^{\rm eq}/n_{\tilde{\gamma}}^{\rm eq})\\&=\Gamma_{R^0\rightarrow\tilde{\gamma}\pi}\left(\frac{M}{m}\right)^{3/2}\exp[-(M-m)/T]\\&=2.0\times10^{-14}r^{3/2}\mathscr{F}(r)\exp[-(r-1)x]\\&\quad\times[\mu_8^5\mu_S^{-4}B]\ \ \mathrm{GeV}.\end{aligned} \tag{25}$$

$\tilde{\gamma}\pi\rightarrow R^0\pi$ *(photino–R^0 conversion).* It is easiest to obtain this term by first considering the term in the equation for $\dot{n}_{R^0}$ due to the reverse process and then using detailed balance:

$$\dot{n}_{R^0}\supset-n_{\pi}n_{R^0}\langle|v|\sigma_{R^0\pi}\rangle. \tag{26}$$

Since the photino–R^0 conversion process creates a $\tilde{\gamma}$ there is a similar term in $\dot{n}_{\tilde{\gamma}}$ with the opposite sign. Now we can write this in a form to calculate the rate for $\tilde{\gamma}$ annihilation by

$$\dot{n}_{\tilde{\gamma}}\supset-\dot{n}_{R^0}=n_{\pi}n_{R^0}\langle|v|\sigma_{R^0\pi}\rangle=\left[\frac{n_{\pi}n_{R^0}}{n_{\tilde{\gamma}}}\langle|v|\sigma_{R^0\pi}\rangle\right]n_{\tilde{\gamma}}. \tag{27}$$

Assuming equilibrium as before, the rate keeping the $\tilde{\gamma}$ in equilibrium can be expressed as

$$\begin{aligned}\Gamma(\tilde{\gamma}\pi\rightarrow R^0\pi)&=\frac{n_{R^0}^{\rm eq}}{n_{\tilde{\gamma}}^{\rm eq}}\,n_{\pi}^{\rm eq}\langle|v|\sigma_{R^0\pi}\rangle\\&=2.7\times10^{-12}r^{5/2}x^{-3/2}\exp(-0.175\mu_8^{-1}x)\\&\quad\times\exp[-(r-1)x][\mu_8^{7/2}\mu_S^{-4}C].\end{aligned} \tag{28}$$

Of course it is the ratio of the reaction rates to the expansion rate that will be used to estimate photino freeze-out. These ratios are given in Table II.

There are two striking features apparent when comparing the magnitudes of the equilibrium reaction rates in Table II. The first feature is that the numerical factor in R^0 self-anihilation is enormous in comparison to the other numerical factors. This simply reflects the fact that R^0 annihilation proceeds through a strong process, while the other processes are all suppressed by a factor of $M_{\tilde{S}}^{-4}$.

The other important feature is the exponential factors of the rates. They will largely determine when the photino will decouple, so it is worthwhile to examine them in detail.

The exponential factor in $\tilde{\gamma}$ self-annihilation is simply $e^{-m/T}$, which arises from the equilibrium abundance of the $\tilde{\gamma}$. It is simple to understand: the probability of one $\tilde{\gamma}$ to find another $\tilde{\gamma}$ with which to annihilate is proportional to the photino density, which contains a factor of $e^{-m/T}$ in the NR limit.

The similar exponential factor in R^0 self-annihilation is also easy to understand. An R^0 must find another R^0 to anihilate, and that probability is proportional to $e^{-M/T}=e^{-rx}$.

Coannihilation is also an exothermic process, so the only exponential suppression is the probability of a $\tilde{\gamma}$ locating the R^0 for coannihilation, proportional to the equilibrium number density of R^0, in turn proportional to $e^{-M/T}=e^{-rx}$.

In inverse decay the exponential factor is $e^{-(r-1)x}=e^{-(M-m)/T}$. The number density of target pions is $e^{-m_\pi/T}$, so this factor is present. It is necessary for the π-$\tilde{\gamma}$ collision to have sufficient center-of-mass energy to create the R^0. This introduces an addition suppression of $e^{-(M-m-m_\pi)/T}$. Combining the two exponential factors gives the result in Table II.

Finally, pion catalysis of photino–gluino conversion involves two exponential suppression factors. The first, $e^{-m_\pi/T}=e^{-0.175\mu_8^{-1}x}$ represents the suppression in the pion number density,[12] and since the mass of the R^0 is greater than the mass of the $\tilde{\gamma}$, there is an additional $e^{-(M-m)/T}$ suppression.

The factors of x and r originate from three places: a factor of x^2 comes from dividing the rates by H, factors of r and x arise for preexponential factors in the number density, and finally they may appear explicitly in the cross section or decay width.

The equilibrium reaction rates divided by H are shown in Figs. 1 and 2 for $r=1.25$, 1.5, 1.75, and 2. In the figures we have assumed $\mu_8=\mu_S=A=B=C=1$. Using the information in Table II it is possible to scale the curves for other values of the parameters.

V. ANALYSIS

Rather than integrate a complete reaction network for the evolution and freeze-out of the photinos, we will assume that

[12]At the temperatures of interest for decoupling, pions might be cheap, but they are not free.

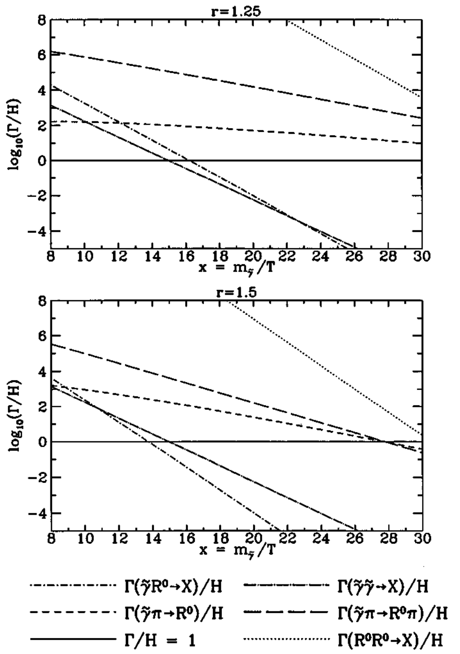

FIG. 1. Equilibrium reaction rates divided by H for $r=1.25$ and 1.5, assuming $\mu_8=\mu_S=1$, and that the factors $A=B=C=1$. The rates can be easily scaled for other choices of the parameters.

the photinos remain in equilibrium so long as there is a reaction depleting the $\tilde\gamma$ abundance that is larger than H. We will then assume that as soon as the rate of the last such reaction drops below H, the photinos immediately freeze out, and the photino-to-entropy ratio is frozen at that value. We will call this approximation the ‘‘sudden’’ approximation.

We can get some idea of the accuracy of the sudden approximation by considering a simple system involving only photino self-annihilation. As discussed in Sec. II, there is a well-developed formalism for calculating the self-annihilation freeze-out of a NR species [11]. Using that formalism in Sec. II, Eq. (8) we find $Y_\infty \simeq 7.4\times 10^{-7}$.

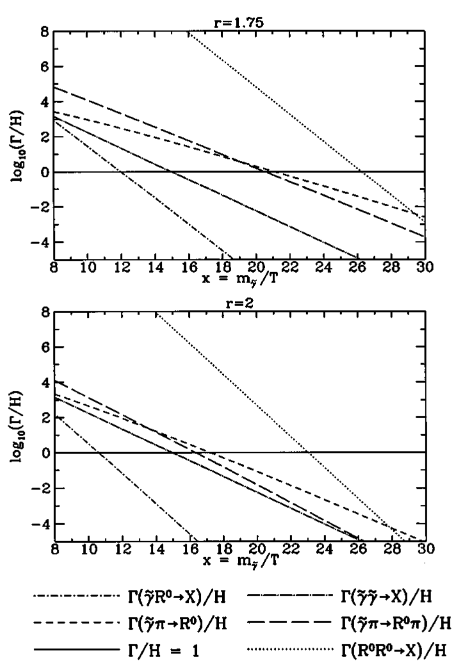

FIG. 2. The same as Fig. 1, but for $r=1.75$ and 2.

Now let us compute Y_∞ using the sudden approximation. From Fig. 1 or Fig. 2, we see that $\Gamma(\tilde\gamma\tilde\gamma\to X)=H$ at $x=14.7$, independent of r. We will denote by x_* the value of x when $\Gamma=H$. Using the sudden approximation that the $\tilde\gamma$ is in LTE until $x=x_*$ and immediately freezes out would give a photino to entropy ratio of (again using 2 degrees of freedom and $g_*=10$)

$$Y_\infty = Y_{\rm eq}(x_*) = 0.145(2/10)x_*^{3/2}\exp(-x_*)$$

$$= 7\times 10^{-7} \quad (\text{using } x_*=14.7). \tag{29}$$

The agreement between Y_∞ obtained using the sudden approximation, Eq. (29), and the usual freeze-out calculation, Eq. (8), suggests that the sudden approximation is a reasonable point of departure for a first look at this phenomenon. Note, however, that the accuracy of the sudden approximation when self-annihilation is the principal photino equilibration mechanism does not guarantee that it is an equally good approximation when interconversion is the important process. The Boltzmann equation when photino self-annihilation dominates can be written [11]

$$\frac{dY}{dx} = \frac{-x\langle\sigma_{\tilde\gamma\tilde\gamma}|v|\rangle s}{H(m)}\,(Y^2 - Y_{\rm eq}^2), \tag{30}$$

where $Y_{\rm eq}(x)$ has the form given in Eq. (29) and $H(m)=1.67g_*^{1/2}m^2/m_{\rm Pl}$. This is to be contrasted with the analogous expression when interconversion dominates:

$$\frac{dY}{dx} = \frac{-x\langle\sigma_{\tilde\gamma\pi\to R^0\pi}|v|\rangle s}{H(m)}\,(Y - Y_{\rm eq})Y_\pi . \tag{31}$$

Here, Y_π is the equilibrium pion to entropy ratio:

$$Y_\pi(x) = 0.145(3/2)(2/10)(r_\pi)x^{3/2}\exp(-r_\pi x). \tag{32}$$

We have introduced $r_\pi \equiv m_\pi/m = 0.175\mu_8^{-1}$, and included the factor (3/2) because the pion has three flavor$\times$spin degrees of freedom in comparison to the photino's two. The

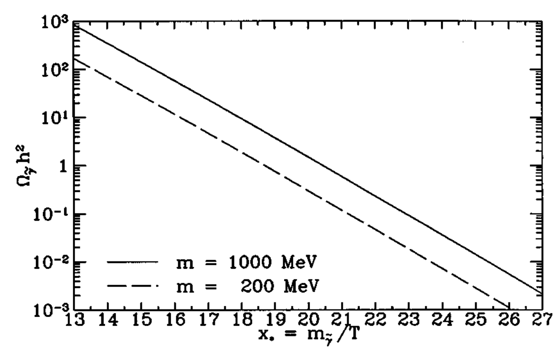

FIG. 3. $\Omega_{\tilde\gamma}h^2$ as a function of x_* assuming the photino stays in equilibrium until x_* and immediately decouples (the sudden approximation).

TABLE III. The value of $\Omega_{\tilde{\gamma}}h^2$ assuming freeze-out at $x=x_*$. $\Omega_{\tilde{\gamma}}h^2=1$ occurs around $x_*=20$, and $\Omega_{\tilde{\gamma}}h^2=10^{-2}$ around $x_*=25$. In the table we have taken $\mu_8=1$.

x_*	12	14	16	18	20	21	22	24	25	26	28	30
$\Omega_{\tilde{\gamma}}h^2$	1660	283	47	7.6	1.2	0.5	0.2	0.03	0.01	0.004	0.0007	10^{-6}

difference in these forms, in particular the much weaker exponential dependence on x for Y_π compared to $Y_{\rm eq}$, is largely responsible for the shallower slope of the interconversion and inverse-decay curves as compared to the self-annihilation curves in Figs. 1 and 2. This shallower slope means that the quality of the sudden approximation in this case is inferior to the self-annihilation case, but probably not significantly in comparison to the large uncertainty due to our present poor knowledge of the cross sections. Closer examination of this question is in progress [20].

Now we proceed using the sudden approximation. Given x_*, we wish to determine $\Omega_{\tilde{\gamma}}h^2$. It is, of course, a very sensitive function of x_*: $\Omega_{\tilde{\gamma}}h^2=2.25\times10^8\,[\mu_8]Y_\infty=6.5\times10^6\,[\mu_8]x_*^{3/2}\exp(-x_*)$. The dependence of $\Omega_{\tilde{\gamma}}h^2$ upon x_* is shown in graphical form in Fig. 3, with specific values presented in Table III.

Since the age of the universe restricts $\Omega_{\tilde{\gamma}}h^2$ to be smaller than unity, x_* must be larger than 20. In order for the relic photinos to be dynamically interesting in structure evolution $\Omega_{\tilde{\gamma}}h^2$ must be larger than 10^{-2}, which obtains for $x_*\lesssim25$. Photinos would dominate the mass of the Universe if $\Omega_{\tilde{\gamma}}h^2\gtrsim0.03$,[13] which would result if $x_*=24$. For $\Omega_{\tilde{\gamma}}=1$ and $h\sim1/2$, x_* must be about 22. Thus we can summarize interesting values of x_* by

$$\begin{array}{rll} x_*\lesssim25; & \Omega_{\tilde{\gamma}}h^2\gtrsim10^{-2}; & \text{dynamically interesting role for photinos,}\\ x_*\lesssim24; & \Omega_{\tilde{\gamma}}h^2\gtrsim0.03; & \text{photinos dominate baryons,}\\ 20\lesssim x_*\lesssim23; & \Omega_{\tilde{\gamma}}h^2\sim0.9; & \text{photinos are the dark matter and } \Omega_{\rm tot=1},\\ x_*\lesssim20; & \Omega_{\tilde{\gamma}}h^2\gtrsim1; & \text{disallowed by age arguments.}\end{array} \tag{33}$$

Now in turn, x_* is exponentially sensitive to $r=M/m$, so limits to the contribution to the present from $\tilde{\gamma}$ will be a sensitive probe of r.

From Figs. 2 and 3, we see that for the canonical choices $\mu_8=\mu_S=A=B=C=1$, either the interconversion process or decay-inverse decay is the last photino reaction to be of importance. It is impossible to say which one because of the uncertainties in the computation of the cross section and the decay width, so we shall consider both possibilities in turn.

If interconversion determines the relic abundance and we make the sudden approximation then we can determine $\Omega_{\tilde{\gamma}}h^2$ as a function of the unknown parameters. Such a graph is given in Fig. 4. From the graph we see that $\Omega_{\tilde{\gamma}}h^2<1$ can result for $r=2.2$ if we allow $\mu_8^2\mu_S^{-4}C$ to be as large as 10^2. We also see that a dynamically interesting value of $\Omega_{\tilde{\gamma}}h^2$ can result for r as small as 1.2 if $\mu_8^2\mu_S^{-4}C=10^{-2}$, although if the interconversion rate is suppressed this much, it is likely inverse decay would govern freeze-out.

A similar calculation can be made assuming that inverse decay is the last operative reaction depleting the photinos. The result of such an analysis is shown in Fig. 5. For $r\gtrsim1.4$ the behavior of the curves is similar to those in Fig. 5, but for small r the effect of phase-space suppression becomes important.

In either case, the conclusion is that for r as large as 2.2, it is possible to have $\Omega_{\tilde{\gamma}}h^2\lesssim1$; with our "central" choice of parameters, $[\mu_8^3\mu_S^{-4}B]=[\mu_8^{3/2}\mu_S^{-4}C]=1$, r must be less than 1.8 in order for $\Omega_{\tilde{\gamma}}h^2\lesssim1$. A value of r as small as 1.2 may result in $\Omega_{\tilde{\gamma}}h^2\gtrsim10^{-2}$; again with the central choice of parameters the limit is $r\gtrsim1.6$.

Although it apparently is not important for realistic parameters, we mention a possible special role for the S^0, $uds\tilde{g}$, the lightest baryon containing a gluino. Since the S^0 has a nonzero baryon number, its abundance is *not* given by Eq. (5) at low temperature because of the nonzero baryon number of the Universe. So long as the strong interactions are maintaining equilibrium between nucleons and S^0's, its abundance should be $n_{S^0}\sim n_N\exp[-(M_{S^0}-m_N)/T]$, where n_N is the nucleon abundance and m_N is the nucleon mass. Thus at very low temperature its abundance will be larger than the R^0 abundance, so the coannihilation and interconversion processes $\tilde{\gamma}S^0\to KN$ and $\tilde{\gamma}N\to KS^0$ are a potential sink of $\tilde{\gamma}$'s which in principle could help keep the $\tilde{\gamma}$ in equilibrium. However, for realistic cross sections, this does not seem to be important at the relevant temperatures. Likewise, although at low enough temperatures there are more nucleons than pions so that $\Gamma_{\tilde{\gamma}N\to R^0N}$ is larger than $\Gamma_{\tilde{\gamma}\pi\to R^0\pi}$, freeze-out has already occurred before the number density of nucleons begins to dominate that of pions.

VI. SUMMARY AND CONCLUSIONS

We have studied the reactions important for the decoupling and freeze-out of photinos having mass m less than about 1.5 GeV. We have found that it is crucial to include the interactions of the photino with the R^0, the gluon-gluino

[13]Nucleosynthesis bounds the contribution from baryons to be about $\Omega_Bh^2\lesssim0.03$.

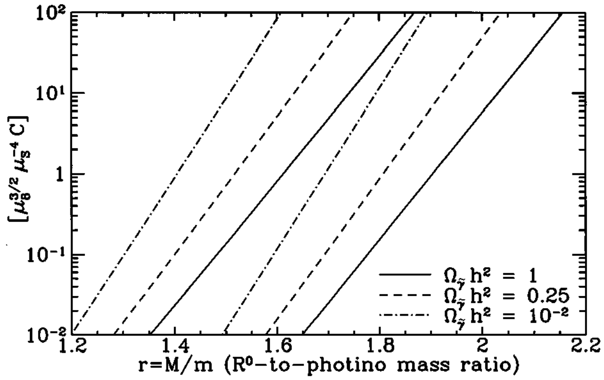

FIG. 4. Assuming $\tilde{\gamma}$ freeze-out is determined by $\tilde{\gamma}$–R^0 conversion, the figure shows as a function of r the values of $[\mu_8^2\mu_S^{-4}C]$ required to give the indicated values of $\Omega_{\tilde{\gamma}}h^2$. The uncertainty band is generated by allowing μ_8 to vary independently over the range $0.5\leqslant\mu_8\leqslant 2$.

bound state whose mass M is expected to lie in the range 1 to 2 GeV. The R^0 has strong interactions and thus annihilates extremely efficiently and stays in thermal equilibrium to much lower temperatures. In this circumstance, photino freeze-out occurs when the rate of reactions converting photinos to R^0's falls below the expansion rate of the Universe. The rate of the $\tilde{\gamma}$–R^0 interconversion interactions which keeps photinos in thermal equilibrium ($\tilde{\gamma}\pi\leftrightarrow R^0\pi$) or R^0 decay-inverse decay ($\tilde{\gamma}\pi\leftrightarrow R^0$), depends on the densities of photinos and pions, rather than on the square of the photino density, as is the case for the self-annihilation process. For photinos of the relevant mass range ($m\sim 800$ MeV), the pion abundance is enormous compared to the photino abundance. Therefore the photinos stay in equilibrium to much higher values of $x\equiv m/T$ than they would if self-annihilation were the only operative process, resulting in a smaller relic density for a given photino mass and cross section. We find using the sudden approximation that light photinos are cosmologically acceptable for a range of $1.2\lesssim r\equiv M/m\lesssim 2.2$. Within this range, if $1.6\lesssim r\lesssim 2$, the photinos are an excellent dark matter candidate. The precise range of r for which the photino accounts for the cold dark matter may shift when the sudden approximation is improved and cross sections are better known. However, the general conclusion is robust: light photinos can account for the dark matter of the Universe for a suitable value of r, which is consistent with theoretical predictions in an attractive class of SUSY-breaking mechanisms [2].

Since $\tilde{\gamma}$–R^0 interconversion governs freeze-out, the usual relation between Ωh^2 and the relic's annihilation cross section [21] is not valid. If inverse decay is the operative process, then there is no direct prediction for the $\tilde{\gamma}$ scattering cross section on matter.[14] If $\tilde{\gamma}\pi\leftrightarrow R^0\pi$ is the operative process, a quantitative solution of the Boltzmann equations can be used to infer properties of this cross section. The cross section of the process relevant for relic detection, $\tilde{\gamma}N\leftrightarrow R^0N$, should be of the same order of magnitude. It will be significantly smaller, more or less by a factor $n_{\tilde{\gamma}}(x_*)/n_\pi(x_*)$, than the conventional cross section used in planning relic detection experiments.

Direct detection of low-mass relic photinos is more difficult than detection of high-mass (say $m\sim 50$ GeV) photinos. In addition the low cross section mentioned above, the average energy deposition is $\langle E\rangle=m^2M_T\langle v^2\rangle/(m+M_T)^2$ where M_T is the target mass. Thus existing and planned experiments using relatively heavy targets are not well adapted to this search. On the positive side, our photino is more likely to have spin-independent couplings to nucleons than expected in the conventional picture [21]. This is because in the SUSY-breaking mechanism which leads to the light photino and gluino under discussion here, the off-diagonal terms in the squark mass-squared matrix can be comparable to the diagonal terms.[15]

Indirect detection via annihilation of gravitationally concentrated photinos [15,22], for instance trapped in the Sun, is unlikely. Because they are low-mass WIMP's, evaporation is much more efficient than in the high-mass case, and they do not concentrate sufficiently. (And, of course, the cross section is smaller than conventionally supposed.)

We also note that if the S^0 is stable, there will be a relic abundance of them, with an abundance relative to baryons determined by M_{S^0} and $x_S\equiv M_{S^0}/T_S$, where T_S is the temperature of S^0 freeze-out. The S^0 mass is expected to be 1.5–2 GeV, so let us define $M_{S^0}=1.5\mu_{1.5}$ GeV. Then

$$\frac{n_S}{n_B}=\frac{1}{4}\left(\frac{M_{S^0}}{m_N}\right)^{3/2}\exp\left[-\frac{M_{S^0}-m_N}{T}\right]$$
$$=\frac{1}{4}\left(\frac{1.5\mu_{1.5}\ \text{GeV}}{0.94\ \text{GeV}}\right)^{3/2}\exp[-(1-0.6/\mu_{1.5})x_S], \qquad (34)$$

where the factor $\frac{1}{4}$ accounts for the fact that the S^0 is a spin-zero state and comes in just one flavor, whereas there are 4 spin$\times$flavor degrees of freedom for the baryons. The S^0 self-annihilation cross section should be comparable to that of the R^0, so ignoring the difference between R^0 and S^0 masses, $x_S\sim rx_{RR}$, where x_{RR} is the value of x at which $\Gamma(R^0R^0\rightarrow X)/H=1$. From Figs. 1 and 2 we see that $rx_{RR}\sim 45$, giving $n_s/n_B\sim 7\times 10^{-9}$ for $\mu_{1.5}=1$, and smaller for larger $\mu_{1.5.}$ Since the S^0's are strongly interacting, even this small an abundance may be detectable. They will be more gravitationally concentrated than standard WIMP's of comparable mass because they dissipate energy through their strong interactions, although they do not form atoms or bind to nuclei.[16]

What, then, is the strategy for testing the proposal that photinos with mass less than or about 1 GeV constitute the

[14] Since the short-distance dynamics entering the matrix element for $R^0\rightarrow\tilde{\gamma}\pi$ is the same as for the scattering reaction $\tilde{\gamma}N\rightarrow R^0N$, these could in principle be related. At this time, however, we do not have sufficient control of the hadron physics involved to make a quantitatively accurate theoretical prediction of the cross sections from the R^0 lifetime.

[15] See Ref. [2] for allowed ranges of the parameters determining the squark mass-squared matrix, μ, $\tan\beta$, and M_0.

[16] If they were stable and could bind to nuclei, they would have been detected in rare isotope searches [1], so that possibility is excluded.

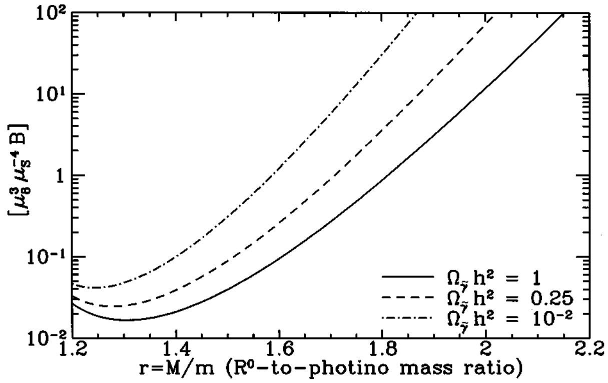

FIG. 5. Assuming $\tilde{\gamma}$ freeze-out is determined by decay-inverse decay, the figure shows as a function of r the values of $[\mu_8^4 \mu_S^{-4} B]$ required to give the indicated values of $\Omega_{\tilde{\gamma}} h^2$.

cold dark matter of the Universe? Of course if an R^0 in the 1 to 2 GeV range could be excluded by laboratory searches, our suggestion for the dark matter would also be excluded. Assuming, though, that these particles are discovered, knowledge of experimentally accessible properties of the photino and R^0 (in particular, their masses, the R^0 lifetime, and the cross section for $R^0 N \rightarrow \tilde{\gamma} N$) coupled with detailed numerical analysis of the freeze-out process, will allow a much more accurate prediction of the relic abundance than has been possible here. Since the relic density is exponentially dependent on r, which will one day be well determined, an accurate quantitative test of this idea will eventually be possible.

In the meantime, theoretical work can elucidate the viability of this proposal. In the class of SUSY-breaking mechanisms relevant to this scenario, the parameters μ, tan β, and M_0, which will determine the photino and gluino masses, are highly constrained [2]. For a specific model and parameter choice, more accurate predictions for the photino and gluino masses can be made. With use of lattice gauge theory, it should be possible to compute the R^0 mass corresponding to a given gluino mass, and thus to determine r for a given model. Lattice gauge calculations could also in principle determine the masses of the other R hadrons and the hadronic matrix elements for the $R^0 - \tilde{\gamma}$ interconversion reactions. For instance, knowledge of the mass of the R_π would allow one to better model $\sigma_{\tilde{\gamma}\pi}$ for given squark masses. With more accurately fixed inputs, a full numerical solution of the coupled Boltzmann equations would be justified.

Therefore, the most important next steps are the following.

(1) Look hard for R hadrons and other new particles predicted by this scenario. Planned kaon experiments may be able to establish evidence for the R^0, and possibly measure its lifetime and mass, as well as the mass of the photino [2].

(2) Do a better job fixing the parameters of the underlying theory, as well as calculating the photino mass produced through radiative corrections.

(3) Use lattice gauge theory to calculate the R^0 mass and check other predictions of this scenario such as the origin of the η' mass [2].

(4) A more complete treatment modeling the photino freeze-out is necessary [20]. An immediate question to address is the quality of the sudden approximation used here. When interconversion is the dominant process, the equation governing the evolution of the photino density has a somewhat different form than in the self-annihilation case, for which the quality of the sudden approximation is well established.

(5) Obtain detailed predictions for the low energy $\tilde{\gamma}$-nucleus cross sections expected in this scenario, and find effective detection techniques for light photinos.

At the very least we have shown that until the value of r is demonstrated to be larger than about 2.2, light photinos are cosmologically acceptable. At best, we have described the scenario for the production and survival of the dark matter of the Universe.

While there is no shortage of candidates for relic dark matter particle species, this proposal extends the idea that photinos may be the dark matter to a previously excluded mass range by incorporating new reactions that determine the photino relic abundance. If this scenario is correct, direct and indirect detection of dark matter might be even more difficult than anticipated. However, the scenario requires the existence of low-mass hadrons, which can be produced and detected at accelerators of moderate energy. Thus particle physics experiments will either disprove this scenario, or make light photinos the leading candidate for dark matter.

ACKNOWLEDGMENTS

G.R.F. was supported in part by the NSF (NSF-PHY-91-21039). E.W.K. was supported by the DOE and NASA under Grant No. NAG5-2788.

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