Observable monochromatic photons from cosmic photino annihilation

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(Received 3 November 1987)

A recent suggestion that the observations of monochromatic photons from annihilation of photinos into quarkonium plus a photon could be a signature of dark-matter photinos is extended to include all S- and P-wave bound quarkonium states for a general interaction Lagrangian and an estimate of the cross sections is given. We also propose the process $\lambda\bar{\lambda}\rightarrow\gamma\gamma$ as a potentially rich source of monochromatic photons and estimate its strength.

I. INTRODUCTION

The dark-matter problem in the Universe has recently attracted attention among astrophysicists and particle physicists.[1] Much work has been devoted to the task of finding experimental signatures for the various candidates for dark matter—axions, photinos, Higgsinos, heavy neutrinos, etc.[2] In the case of photinos (λ) it has recently been suggested by Srednicki, Theisen, and Silk[3] that the annihilation process $\lambda\bar{\lambda}\rightarrow V\gamma$ might give a detectable photon signal (V is the vector-meson bound state of heavy quarks, e.g., $c\bar{c}$). This analysis, which was based on a pointlike coupling between the $Q\bar{Q}$ pair and the vector meson, has been criticized by Rudaz[4] who found a lower branching ratio for $\lambda\bar{\lambda}\rightarrow V\gamma$ using a more realistic bound-state description of the vector meson. For $m_\lambda<4$ GeV, photons from $\lambda\bar{\lambda}\rightarrow J/\psi\gamma$ should still be detectable using one of the high-resolution detectors that have been proposed.[5] High-energy resolution enables taking advantage of the near monochromaticity of the lines which is due to the low (galactic) relative velocities of the annihilating λ pair ($v_{\rm rel}/c=10^{-3}$).

In this paper we extend the analysis of Rudaz to include the processes $\lambda\bar{\lambda}\rightarrow({}^{2S+1}L_J)+\gamma$ for all bound $Q\bar{Q}$ P- and S-wave states and for a more general interaction Lagrangian. In addition, we propose the process $\lambda\bar{\lambda}\rightarrow\gamma\gamma$ as a potentially rich source of monochromatic photons and make an estimate of its strength.

II. AMPLITUDES FOR RADIATIVE QUARKONIUM PRODUCTION

We write the effective interaction Lagrangian for the interaction between a λ pair and a fermion (f) pair as

$$L_{\rm eff}=l_\mu[\bar{f}\gamma^\mu(v_f+\gamma^5a_f)f]\ , \tag{1}$$

where $l_\mu=\bar{\lambda}\gamma_\mu(v_\lambda+a_\lambda\gamma_5)\lambda$. When f is a quark field, a sum over colors is implied. There is an important distinction between the case when λ is a Majorana fermion (such as a photino) and in the straightforward case when λ is a Dirac fermion (such as a heavy neutrino). In the Majorana case, when passing from Eq. (1) to the Feynman rules a_λ effectively gets multiplied by a factor of 2 whereas v_λ vanishes. We will write the various amplitudes as if λ is a Dirac fermion, but will take the last remark into account when presenting cross sections for photinos. For supersymmetric particles such as photinos, Higgsinos, etc., the effective couplings in (1) have been expressed in terms of gauge couplings and scalar superpartner masses by Ellis *et al.*[2]

In models where the scalars associated with the left- and right-handed fermions, respectively, are nearly degenerate in mass the coupling v_f is small whereas it may be important for large mass differences [generally $v_f/a_f\approx(\tilde{m}_L^2-\tilde{m}_R^2)/(\tilde{m}_L^2+\tilde{m}_R^2)$].

With the effective Lagrangian (1) governing the coupling specifically between $\lambda\bar{\lambda}$ and a heavy-quark pair, we must prescribe the bound-state model before we can calculate the process $\lambda\bar{\lambda}\rightarrow({}^{2S+1}L_J)+\gamma$ depicted in Fig. 1. We adopt the point of view of Rudaz[4] that an adequate description of the process is provided by a nonrelativistic

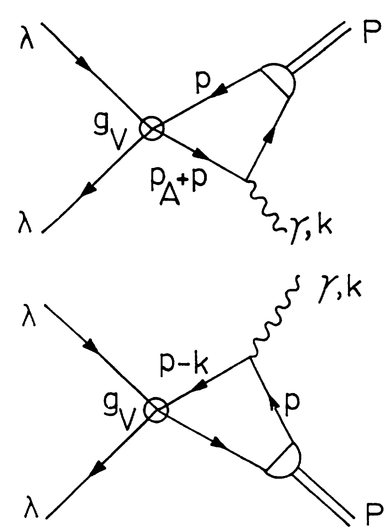

FIG. 1. Diagrams contributing to the process $\lambda\bar{\lambda}\rightarrow P\gamma$, where P is a ${}^{2S+1}L_J$ quarkonium bound state.

bound-state model which has proven to be remarkably successful in the calculation of charmonium and bottomonium properties.

It is convenient to use the formalism developed in Ref. 6. Working in the rest frame of the produced heavy meson one finds

$$M(\lambda\bar{\lambda}\rightarrow {}^{2S+1}L_J+\gamma)=\frac{ie\sqrt{N_C}e_Q}{4m_P\sqrt{\pi m_P}}\int\frac{d^3p}{(2\pi)^3}\tilde{R}_L(p)\frac{m_P^2\mathrm{Tr}(X_P\tilde{A})}{2E(E+m_Q)}, \quad (2)$$

where the positive-energy amplitude $\tilde{A}$ is given by

$$\tilde{A}=(1+\not{P}/m_P)(\not{p}_-+m_Q)\times\left[\not{\epsilon}^*(k)\frac{1}{\not{p}_-+\not{k}-m_Q}\not{l}(v_Q+a_Q\gamma^5)+\not{l}(v_Q+a_Q\gamma^5)\frac{1}{\not{p}_--\not{p}_A-m_Q}\not{\epsilon}^*(k)\right](\not{p}_+-m_Q). \quad (3)$$

Here P,p_-,p_+,k,p_A are the momenta of the meson, quark, antiquark, photon, and the fermion current l_μ, respectively. The photon polarization vector is denoted by ϵ; m_P and m_Q are the meson and quark masses.

In (2), e_Q is the quark charge in units of the proton charge, and $\tilde{R}_L(p)$ is the radial part of the L-wave function in momentum space. The various state projectors X_P are given by

$$0^{-+}:\ X_P=\gamma^5, \quad (4a)$$

$$1^{--};\ X_P=-\not{\epsilon}(\lambda), \quad (4b)$$

$$1^{+-}:\ X_P=\sqrt{3}\vec{\epsilon}(\lambda)\cdot\hat{\mathbf{p}}\gamma^5, \quad (4c)$$

$$0^{++}:\ X_P=(\not{p}_--\not{p}_+)/(2p), \quad (4d)$$

$$1^{++}:\ X_P=i\sqrt{3}\epsilon_{\sigma\alpha\beta\rho}P^\alpha\epsilon^\beta(\lambda)\times(p_--p_+)^\sigma\gamma^\rho/(2\sqrt{2}pm_P), \quad (4e)$$

$$2^{++}:\ X_P=(\sqrt{3}/2p)S_{\rho\sigma}(\lambda)(p_--p_+)^\rho\gamma^\sigma. \quad (4f)$$

Here $\epsilon(\lambda)$ and $S(\lambda)$ are polarization four-vectors and four-tensors for spin-1 and spin-2 states, respectively, and $\hat{\mathbf{p}}$ is a unit three-vector in momentum space. Inserting these expressions into (3), performing the indicated traces and retaining only lowest-order nonvanishing terms in the momentum integration, one obtains in the small-binding limit $(m_Q\rightarrow m_P/2)$ the following total amplitudes:

$$0^{-+}:\ M=F_1\epsilon_\nu^*(k)l_\mu P_\rho k_\sigma\epsilon^{\rho\sigma\mu\nu}, \quad (5a)$$

with

$$F_1=2iee_Q\sqrt{N_C}v_Q\psi_S(0)/[\sqrt{2m_Q}(m_\lambda^2-m_Q^2)]; \quad (5b)$$

$$1^{--}:\ M=F_2\epsilon_\nu^*(k)l_\mu\epsilon_\rho(\lambda)k_\sigma\epsilon^{\rho\sigma\mu\nu}, \quad (6a)$$

with

$$F_2=2iee_Q\sqrt{N_C}a_Q\sqrt{2m_Q}\psi_S(0)/(m_\lambda^2-m_Q^2); \quad (6b)$$

$$1^{+-}:\ M=F_3\epsilon_\nu^*(k)l_\mu\epsilon_\rho^*(\lambda)\times\{k^\rho[-g^{\mu\nu}+k^\mu P^\nu/(k\cdot P)]+P^\mu[-g^{\rho\nu}+k^\rho P^\nu/(k\cdot P)]\}, \quad (7a)$$

with

$$F_3=2i\sqrt{3}\sqrt{N_C}ee_Qa_QR_P'(0)(m_\lambda^2-m_Q^2)^{-1}/\sqrt{\pi m_P}; \quad (7b)$$

$$0^{++}:\ M=F_4\epsilon_\nu^*(k)l_\mu[g^{\mu\nu}-P^\nu k^\mu/(k\cdot P)], \quad (8a)$$

with

$$F_4=\frac{i\sqrt{2m_Q}\sqrt{N_C}ee_Qv_Q}{\sqrt{\pi}(m_\lambda^2-m_Q^2)}\frac{m_\lambda^2-3m_Q^2}{m_Q^2}R_P'(0); \quad (8b)$$

$$1^{++}:\ M=F_5\epsilon_\nu^*(k)l_\mu\epsilon_\rho^*(\lambda)\times(m_\lambda^2\epsilon_{\rho\sigma\mu\nu}k^\sigma+\epsilon_{\nu\sigma\rho\gamma}P^\sigma k^\gamma p_A^\mu/4)/(m_\lambda^2-m_Q^2), \quad (9a)$$

with

$$F_5=-2\sqrt{6}\sqrt{N_C}ee_Qv_QR_P'(0)/[\sqrt{2\pi m_Q}(m_\lambda^2-m_Q^2)]; \quad (9b)$$

$$2^{++}:\ M=F_6\epsilon_\nu^*(k)l_\mu S_{\rho\sigma}^*(P)\times\{k^\sigma(g^{\mu\nu}k^\rho-g^{\nu\rho}k^\mu)-g^{\mu\rho}[k^\sigma P^\nu-(P\cdot k)g^{\sigma\nu}]\}, \quad (10a)$$

with

$$F_6=-i\sqrt{3}\sqrt{N_C}ee_Qv_Q\sqrt{2m_Q}R_P'(0)/[\sqrt{\pi}(m_\lambda^2-m_Q^2)^2]. \quad (10b)$$

From these expressions we see that the vector coupling v_Q contributes to 0^{-+}, 0^{++}, 1^{++}, and 2^{++} production, whereas the axial-vector coupling a_Q contributes to 1^{--} and 1^{+-} production, in accordance with C invariance. For the case of photinos, only the latter two states are then important in the limit of degenerate squark masses.

III. CROSS SECTIONS

Using the interaction Lagrangian (1), the cross section $\lambda\bar{\lambda}\rightarrow f\bar{f}$ is calculated to be (for Dirac λ's)

$$\begin{aligned}\sigma(\lambda\bar{\lambda}\rightarrow f\bar{f})=N_C\beta_Q[12\pi\beta_\lambda(1-\beta_Q^2)]^{-1}[&a_\lambda^2a_Q^2(m_Q^2\beta_\lambda^2\beta_Q^2+6m_\lambda^2\beta_Q^4+3m_Q^2\beta_Q^2+3m_\lambda^2)\\&+a_\lambda^2v_Q^2(m_Q^2\beta_\lambda^2\beta_Q^2-3m_\lambda^2\beta_Q^4+9m_\lambda^2\beta_Q^2-3m_Q^2\beta_Q^2+6m_Q^2-6m_\lambda^2)\\&+v_\lambda^2a_Q^2\beta_Q^2(m_Q^2\beta_\lambda^2-3m_\lambda^2\beta_Q^2+3m_\lambda^2+3m_Q^2)\\&+v_\lambda^2v_Q^2(m_Q^2\beta_\lambda^2\beta_Q^2-3m_\lambda^2\beta_Q^2-3m_Q^2\beta_Q^2+3m_\lambda^2+6m_Q^2)],\end{aligned} \quad (11)$$

where β_Q and β_λ are the quark and λ velocities in the center-of-mass system (c.m.s). For our applications, the limit $\beta_\lambda = v_{\rm rel}/2 \to 0$ is of interest. From (11) one immediately finds

$$(\sigma v_{\rm rel})_{v_{\rm rel}\to 0} = (2\pi)^{-1} N_C \beta_Q \times [a_\lambda^2 a_Q^2 m_Q^2 + 2v_\lambda^2 a_Q^2 (m_\lambda^2 - m_Q^2) + v_\lambda^2 v_Q^2 (2m_\lambda^2 + m_Q^2)] \ . \tag{12}$$

For Majorana λ's, $v_\lambda = 0$ and $a_\lambda \to 2a_\lambda$. Thus,

$$(\sigma v_{\rm rel})^{\rm Maj}_{v_{\rm rel}\to 0} = (2N_C/\pi)\beta_Q a_\lambda^2 a_Q^2 m_Q^2 \ . \tag{13}$$

Sometimes, the expression for $\sigma v_{\rm rel}$ is required to second order in $v_{\rm rel}^2$. This is

$$(\sigma v_{\rm rel})^{\rm Maj} = \pi^{-1} 2 m_\lambda^2 N_C a_\lambda^2 a_Q^2 \sqrt{1-\mu^2} \times (\mu^2 + v_{\rm rel}^2 \{2 + 3\mu^2[-17/3 + (1-\mu^2)^{-1}]/4\}/6) \ , \tag{14}$$

where $\mu = m_Q/m_\lambda$. This formula has often been given incorrectly in the literature. Using now the expressions (5a)–(10a) for the amplitudes for the radiative quarkonium production, one obtains the corresponding cross sections (for Dirac λ's), to lowest order in $v_{\rm rel}$:

$$\sigma(0^{-+}\gamma) = |F_1|^2 m_\lambda^4 v_\lambda^2 (1-\mu^2)^3/(2\pi v_{\rm rel}) \ , \tag{15}$$

$$\sigma(1^{--}\gamma) = |F_2|^2 m_\lambda^2 (1-\mu^2)^3 \times [a_\lambda^2 + v_\lambda^2(1+\mu^{-2})]/(8\pi v_{\rm rel}) \ , \tag{16}$$

$$\sigma(1^{+-}\gamma) = |F_3|^2 m_\lambda^2 (1-\mu^2) \times [a_\lambda^2(1+\mu^2)^2 + v_\lambda^2(\mu^{-2} - 1 - \mu^2 + \mu^4)]/(8\pi v_{\rm rel}) \ , \tag{17}$$

$$\sigma(0^{++}\gamma) = |F_4|^2 v_\lambda^2 (1-\mu^2)/(8\pi v_{\rm rel}) \ , \tag{18}$$

$$\sigma(1^{++}\gamma) = |F_5|^2 m_\lambda^2 v_\lambda^2 (\mu^{-2} - \mu^2)/(8\pi v_{\rm rel}) \ , \tag{19}$$

and

$$\sigma(2^{++}\gamma) = |F_6|^2 m_\lambda^4 v_\lambda^2 (1-\mu^2)^3 \times (\mu^{-4} + 3\mu^{-2} + 6)/(12\pi v_{\rm rel}) \ . \tag{20}$$

To be specific, we now restrict ourselves to the case of photino annihilation. Then $v_\lambda = 0$, $a_\lambda \to 2a_\lambda$ and we see from (15)–(20) that only 1^{--} and 1^{+-} production are nonvanishing. For the 1^{--} case, which is the only one which has previously been considered in the literature, we recover the result of Rudaz:[4]

$$\sigma(\lambda\lambda \to 1^{--} + \gamma)_{\rm photino} = 4e^2 e_Q^2 N_C m_Q |\psi_S(0)|^2 \times a_Q^2 a_\lambda^2 (1-\mu^2)/(\pi m_\lambda^2 v_{\rm rel}) \ , \tag{21}$$

whereas we get for 1^{+-} production

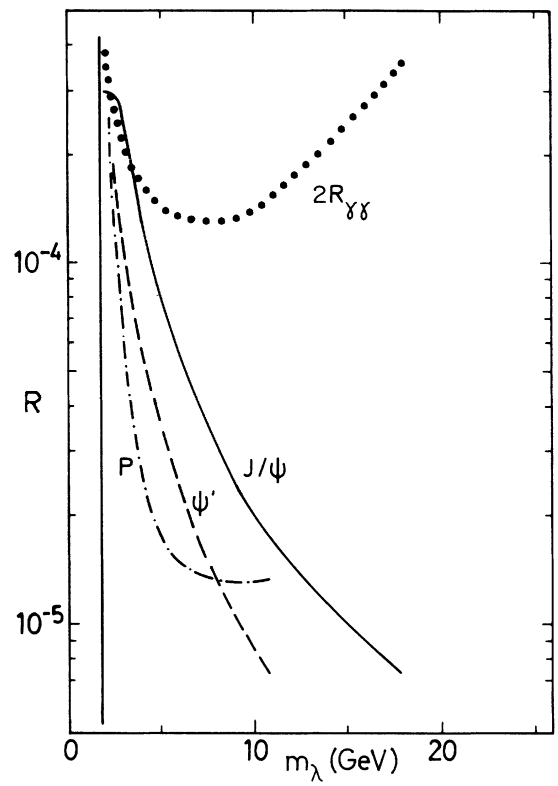

FIG. 2. The ratios $R_V = \sigma(\lambda\bar\lambda \to 1^{--} + \gamma)_{\rm photino}/\sigma(\lambda\bar\lambda \to c\bar c)_{\rm photino}$ ($V = J/\psi, \psi'$) and $R_P = \sigma(\lambda\bar\lambda \to 1^{+-} + \gamma)_{\rm photino}/\sigma(\lambda\bar\lambda \to c\bar c)_{\rm photino}$ ($P = {}^1P_1$) shown as a function of m_λ, the mass of the annihilating photinos. The corresponding ratio for the first excited 1^{--} state is also shown. The dotted curve is the ratio $2R_{\gamma\gamma} = 2\sigma(\lambda\bar\lambda \to \gamma\gamma)/\sigma(\lambda\bar\lambda \to c\bar c)$ as a function of the photino mass.

$$\sigma(\lambda\lambda \to 1^{+-} + \gamma)_{\rm photino} = 3e^2 e_Q^2 N_C (\pi^2 m_Q v_{\rm rel})^{-1} \times a_Q^2 |R_P'(0)|^2 \times a_\lambda^2 \mu^2 (1+\mu^2)^2/(1-\mu^2) \ . \tag{22}$$

A remarkable feature of the result (22) is the factor $1-\mu^2$ in the denominator which enhances this mode compared to 1^{--} production for m_λ values close to m_Q. This is illustrated in Fig. 2, where the ratios of (21) and (22) to the $Q\bar Q$ cross section (13) are shown for charmonium. In Fig. 2 also cross section for the radial excitation ψ' is shown.

The fact that the 1^{+-} photon line for large values of the λ mass is brighter than those from the excited 1^{--} states is important, since to extract the mass of the annihilating λ's one needs at least two observed photon lines which must be related to charmonium states of known mass.

Following the analysis of Refs. 3 and 4, we conclude that these processes may give an observable photon signal if $m_\lambda < 4$ GeV.

IV. TWO-PHOTON ANNIHILATION OF PHOTINOS

In the previous sections we have seen that cosmic photino annihilation into 1^{--} or 1^{+-} charmonium states plus a photon could give a detectable photon signal. We now turn to another, potentially interesting mechanism for production of nearly monochromatic, cosmic photons, namely, photino annihilation into two photons.

The $\lambda\bar{\lambda}\to\gamma\gamma$ process is zero at the tree level but is induced in higher order through loop diagrams. A full calculation consists of evaluating several hundred diagrams. In this paper we will only give a simplified calculation based on several simplifying assumptions which we think at least gives a reliable estimate of the order of magnitude for this cross section.

First, we only keep the contribution from lepton and quark loops (not $W^{\pm}$, etc.). Second, we assume that the supersymmetric spin-zero partners of quarks and leptons are so heavy that the effective interaction (1) can be used when integrating over the fermion loops (Fig. 3). Furthermore, the contribution from fermions circulating the loop can be neglected.

In the initial state, the $\lambda\bar{\lambda}$ pair can be in a 1^{++} or 0^{-+} state. Since a 1^{++} cannot decay into two real photons due to Yang's theorem,[7] we need only consider the pseudoscalar contribution. Using the projector for the pseudoscalar initial state

$$P_{\rm ps}^{\rm in}=m_\lambda\gamma^5(1-\not{P}/2m_\lambda)/\sqrt{2}\ , \tag{23}$$

one gets

$$\mathrm{Tr}(P_{\rm ps}^{\rm in}\gamma_\mu\gamma^5)=-\sqrt{2}P_\mu\ . \tag{24}$$

Using then the divergence condition for the axial-vector current (there is no anomaly, since the loop diagram is convergent when "resolving" the $\lambda\bar{\lambda}Q\bar{Q}$ vertex) $P_\mu j^{\mu 5}=2m_Q j^5$, we find the matrix element

$$M=\epsilon_{1\mu}^*(k_1)\epsilon_{2\nu}^*(k_2)M^{\mu\nu}\ , \tag{25}$$

with

$$M_{\mu\nu}=-16\sqrt{2}m_Q^2 a_\lambda a_Q N_C(-iee_Q)^2\epsilon_{\mu\nu\alpha\beta}k_1^\alpha k_2^\beta I\ , \tag{26}$$

where

$$I=\int\frac{d^4k}{(2\pi)^4[(k-k_1)^2-m_Q^2][(k-k_1-k_2)^2-m_Q^2](k^2-m_Q^2)}$$
$$=-iF(1/\mu_Q^2)/(32\pi^2m_\lambda^2)\ . \tag{27}$$

Here

$$F(x)=\begin{cases}\arcsin^2\sqrt{x}\ , & x<1\ ,\\ [\pi^2-\ln^2(\sqrt{x}+\sqrt{x-1})^2]/4 \\ \quad +i\pi\ln(\sqrt{x}+\sqrt{x-1}), & x>1\ .\end{cases} \tag{28}$$

This gives

$$\sigma(\lambda\bar{\lambda}\to\gamma\gamma)=m_\lambda^2a_\lambda^2\alpha^2v_{\rm rel}^{-1}\pi^{-3}\times\Big|\sum_f\mu_f^2a_fQ_f^2F(1/\mu_f^2)\Big|^2\ , \tag{29}$$

where the sum is over all quarks and leptons (including a factor N_C for color) and a top-quark mass of 50 GeV has been assumed (our results are quite insensitive to this).

To calculate the branching ratio for $\lambda\bar{\lambda}\to\gamma\gamma$ to $\lambda\bar{\lambda}\to c\bar{c}$ we assume a common mass $\tilde{m}$ for all squarks and

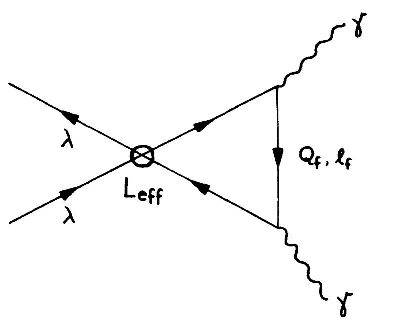

FIG. 3. Effective loop diagrams that contribute to the process $\lambda\bar{\lambda}\to\gamma\gamma$.

sleptons. Then $a_\lambda^2a_f^2=(2\pi\alpha Q_f^2/\tilde{m}^2)^2$, and

$$\sigma(\lambda\bar{\lambda}\to c\bar{c})v_{\rm rel}=128\pi\alpha^2m_c^2\beta_c/(27\tilde{m}^4)\ , \tag{30}$$

whereas

$$\sigma v_{\rm rel}(\gamma\gamma)=4\alpha^4\pi^{-1}m_\lambda^2\tilde{m}^{-4}\times\Big|\sum_f\mu_f^2Q_f^4F(1/\mu_f^2)\Big|^2\ . \tag{31}$$

The ratio $\gamma\gamma/c\bar{c}$ is shown in Fig. 2 as a function of m_λ. This function falls less rapidly than the corresponding ratios for $1^{--}\gamma$ and $1^{+-}\gamma$ production of photons, and in fact rises from $m_\lambda>10$ GeV. It can be seen that in the range of m_λ between 1 and 6 GeV it is typically of the order of a few times 10^{-4}. This makes it potentially interesting as a source of cosmic photons. It should be noted that to give a fair comparison between this process and radiative quarkonium production, we have multiplied the $\gamma\gamma$ branching ratio by a factor of 2 in the figure, since here we get two photons per annihilation.

The important feature of Eq. (29), that the cross section is proportional to m_λ^2, is independent of our approximations and comes just from dimensional arguments. In fact, recalling that slow photinos annihilate from a pseudoscalar state, the relatively large 2γ width is not surprising (cf. $\pi^0\to2\gamma$).

It has recently been suggested that Higgsinos, rather than photinos, constitute the dark matter.[8] Our results for the radiative quarkonium production are unchanged in this case. For the $\gamma\gamma$ process, assuming that Z^0 exchange is the main contribution to the $\bar{\lambda}\lambda\bar{f}f$ vertex, we find roughly a 20% reduction of the $\gamma\gamma$ production ratio. It should be noted, though, that the present indications of a heavy top squark ($m_t > m_W$) could make the t-channel squark exchange $\sim(m_t v_1/\tilde{m} v_2)^2$ (here v_1 and v_2 are the vacuum expectation values of the two Higgs doublets needed in supersymmetric models) substantial and in fact increase our estimates for $\gamma\gamma$ production.

A consequence of our analysis of $\gamma\gamma$ production is also that we expect $\lambda\bar{\lambda}\to gg$ (g is the gluon) to be substantial. (This channel has not been considered in previous work.) The cross section for gg production is simply obtained from Eq. (31) by restricting the fermion sum to quarks and by substituting $N_C^2\alpha^2 Q_f^4 \to 2\alpha_S^2$. As a result, we estimate the branching ratio for the gg channel to around 5–10%.

V. CONCLUSIONS

The results of the present investigation can now be summarized as follows.

(i) Monochromatic photons resulting from the annihilation of Majorana particles can be a signature for such types of dark cosmic matter in the galactic halos, as originally suggested by Srednicki, Theisen, and Silk.[3] For photino annihilation into quarkonium and gamma only the 1^{--} and 1^{+-} bound states contribute. For the 1^{--} case we recover the result by Rudaz.[4] For the 1^{+-} case it turns out that for large values of the photino mass this line is brighter than the 1^{--} line. These process might give rise to observable lines provided the photino mass is <4 GeV.

(ii) Some recently conceived high-resolution gamma-ray spectrometers described in Ref. 5 seem to open up further possibilities for studying discrete cosmic gamma lines.

(iii) An estimate of the branching ratio for $\lambda\bar{\lambda}\to\gamma\gamma$ shows that the ratio $\gamma\gamma/c\bar{c}$ is of the order 10^{-4} for m_λ in the range between 1 and 6 GeV. This makes it potentially interesting as a further source of cosmic photons.

ACKNOWLEDGMENT

We would like to thank P. Carlson for useful discussions and presentation of his own ideas on γ-ray detection and for bringing the proposal of Eichler and Adams to our attention.

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