Eternal annihilations of light photinos

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In a class of low-energy supersymmetry models the photino is a natural dark-matter candidate. We investigate the effects of post-freeze-out photino annihilations that generate electromagnetic cascades and lead to photodestruction of ^{4}He and subsequent overproduction of D and ^{3}He. We also generalize our analysis of electromagnetic showers to include those from a generic dark-matter component whose relic abundance is *not* determined by the self-annihilation cross section. [S0556-2821(96)02618-5]

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I. INTRODUCTION

There are good reasons to consider models of low-energy supersymmetry [1] in which dimension-3 supersymmetry-breaking operators are highly suppressed. The low-energy features and possible new signatures of such an attractive form of supersymmetry breaking have been extensively outlined by Farrar and Masiero [2]. In most of these models photinos $\widetilde{\gamma}$ and gluinos $\widetilde{g}$ are very light, and the lightest R-odd particle[1] may be a color-singlet state containing a gluino, the R^0, with a mass m_{R^0} in the 1 to 2 GeV range.

It has been recently pointed out by Farrar and Kolb [3] that a photino $\widetilde{\gamma}$ slightly lighter than the R^0, in the mass range 100 to 1400 MeV, would survive as the relic R-odd species and might be an attractive dark-matter candidate. Indeed, they found that it is crucial to include the interactions of the photino with the R^0. The R^0 has strong interactions and thus annihilates extremely efficiently and remains in equilibrium to temperatures much lower than its mass. In these circumstances, photino freeze out is *no longer* determined by self-annihilations into fermions pairs $\widetilde{\gamma}\widetilde{\gamma}\leftrightarrow f\overline{f}$ (which would result in an unacceptable photino relic abundance for such a low mass range), but occurs when the rate of reactions converting photinos to R^0's falls below the expansion rate of the Universe. The rate of the $\widetilde{\gamma}\leftrightarrow R^0$ interconversion interactions which keep photinos in equilibrium ($\widetilde{\gamma}\pi\leftrightarrow R^0\pi$ or R^0 $\widetilde{\gamma}\pi\leftrightarrow R^0$), depends upon the densities of photinos and pions, rather than on the square of the photino density, as in the case for the self-annihilation processes. For photinos in the relevant mass range ($m_{\widetilde{\gamma}}\sim 800$ MeV), the pion abundance is enormous compared to the photino abundance. Therefore, the photinos stay in equilibrium to much higher values of $x\equiv m_{\widetilde{\gamma}}/T$ 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. Light photinos with a mass in the range $1.2\lesssim r\equiv m_{R^0}/m_{\widetilde{\gamma}}\lesssim 2.2$ are cosmologically acceptable and in the range $1.6\lesssim r\lesssim 2$ are excellent dark matter candidates.

The aim of the present paper is to investigate whether in the scenario outlined in [3], light photino self-annihilations into fermion pairs have an impact on the successful predictions of standard big-bang nucleosynthesis even though self-annihilations are not relevant in determining the relic photino abundance. The destructive high-energy photons coming from the self-annihilation products of primordial photinos, including $e^{\pm}$, $\mu^{\pm}$, and γ's, may indeed cause photofission of the primordially produced light nuclei [4], upsetting the agreement between big-bang nucleosynthesis and the observed element abundances.

According to the standard lore, the source of destructive high-energy photons may be neglected because self-annihilations (SA's) typically freeze out at a temperature $T_{\rm SA}\sim m_{\widetilde{\gamma}}/15$ [not to be confused with the freeze-out temperature of the photino relic abundance $T_*\sim m_{\widetilde{\gamma}}/(20-25)$ determined by $\widetilde{\gamma}\leftrightarrow R^0$ conversion]. For photino masses in the dangerous range, $m_{\widetilde{\gamma}}$ greater than a few MeV, this freeze out occurs well before the light elements first become vulnerable to photodissociation at temperatures of about a few keV. By this late epoch, any photons and electrons generated before freeze out have harmlessly thermalized. However, after freeze out, occasional self-annihilations still may take place. Although rare on the expansion time scale, residual annihilations continue to produce high-energy photons well after nucleosynthesis ends, thereby placing the survival of the light nuclei in jeopardy [5].

We explore below the consequences of these eternal light-photino self-annihilations by computing the residual annihilation rate and the corresponding photofission yields. The typical feature of the model at hand is that the production rate of dangerous high-energy photons and the relic abundance of light photinos are *not* determined by the same cross section, the former being dependent on the light-photino

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[1] R parity is a multiplicative quantum number, exactly conserved in most supersymmetric models, under which ordinary particles have $R=1$ while their superpartners have $R=-1$.

thermally averaged self-annihilation cross section $\langle \sigma_{SA}|v|\rangle$ and the latter by the $\tilde{\gamma} \leftrightarrow R^0$ conversion rate. We then generalize our analysis to any case in which the relic abundance of a generic dark-matter candidate X is not determined by its self-annihilation cross section, as assumed in Ref. [5], but rather by some other generic processes, a typical case being the presence of a slightly heavier particle X' (e.g., the R^0 in Ref. [3]) with relative interconversion processes $X \leftrightarrow X'$ (usually called coannihilation [6]).

II. ETERNAL ANNIHILATIONS

In the Boltzmann equation [7] for the evolution of the $\tilde{\gamma}$-number density there are several terms, including photino self-annihilation ($\tilde{\gamma}\tilde{\gamma} \leftrightarrow X$), coannihilation ($\tilde{\gamma}R^0 \leftrightarrow X$), inverse decay ($\tilde{\gamma}\pi\pi \leftrightarrow R^0$), and photino-$R^0$ conversion ($\tilde{\gamma}\pi \leftrightarrow R^0\pi$) (see Refs. [3] and [8] for a detailed analysis). In this section we are only interested in the terms in the Boltzmann equation which may provide a source for destructive electromagnetic cascades at low temperatures.

There are two freeze-out times of importance in our problem. The first is the freeze-out time of photino self-annihilation. We will define this time as the time when the reaction rate for $\tilde{\gamma}\tilde{\gamma} \leftrightarrow X$ becomes smaller than the expansion rate, $H = 1.66 g_*^{1/2} T^2/M_{\rm Pl} = 4.4 \times 10^{-19} x^{-2} (m_{\tilde{\gamma}}/\text{GeV})$ GeV, with g_* the relativistic degrees of freedom at temperature T, and $x = m_{\tilde{\gamma}}/T$. We denote this time as t_{SA}. The other freeze-out time is when the *total* rate for photino creation or destruction becomes smaller than H. We denote this time by t_*. In previous considerations of eternal annihilations it was assumed that freeze out is determined by self-annihilation, i.e., $t_{SA} = t_*$. However, as discussed above, in the light-photino scenario developed by Farrar and Kolb the photino remains in equilibrium after t_{SA}, i.e., $t_* > t_{SA}$.

We are interested in the epoch long after freeze out, $t \gg t_*$, when the number density to the entropy density ratio $Y_{\tilde{\gamma}} = n_{\tilde{\gamma}}/s$ [with $s = (2\pi^2/45) g_{*S} T^3$ the entropy density at temperature T] has ceased to decrease significantly, i.e., $Y_{\tilde{\gamma}} \simeq Y_\infty =$ constant for $t \gg t_*$. The asymptotic value of the photino-to-entropy ratio, Y_∞, may always be expressed in terms of the fraction of the critical density contributed by the photino at the present epoch:

$$\Omega_{\tilde{\gamma}}\, h^2 = \frac{\rho_{\tilde{\gamma}}}{\rho_C} = 2.8 \times 10^8 \left(\frac{m_{\tilde{\gamma}}}{\text{GeV}} \right) Y_\infty , \tag{1}$$

where $\rho_C = 1.05 \times 10^{-5} h^2$ GeV cm^{-3} is the present critical density.

Well after the freeze out of the photino abundance, the change in time of the fraction of photino number density contributing to fermion-pair production is then given by

$$\dot{n}_{\tilde{\gamma}} + 3 H n_{\tilde{\gamma}} = -\langle \sigma_{SA}|v|\rangle Y_\infty^2 s^2 . \tag{2}$$

It is easy to check that for the temperature of interest here, $T \ll T_*$, self-annihilation into fermion pairs is the dominant source for photino number density depletion so that Eq. (2) actually describes the change in time of the whole photino number density. Using the expression for Y_∞ from Eq. (1), Eq. (2) will be our starting point in the next sections to analyze the influence of eternal self-annihilation on primordial nucleosynthesis.

III. ELECTROMAGNETIC CASCADES

Let us begin by examining the manner in which annihilations of light photinos into light leptons (electrons or muons) generate electromagnetic cascades in the radiation-dominated thermal plasma of the early Universe [9].

The epoch of interest for cascade nucleosynthesis has temperatures smaller than a few keV, by which time production of the light elements D, He, and Li by big-bang nucleosynthesis has completed. At this epoch, electron-positron pairs are no longer in equilibrium and blackbody photons, γ_{bb}, constitute the densest target for electromagnetic cascade development.

A cascade is initiated by a high-energy lepton (or photon) coming from the self-annihilating photinos and develops rapidly in the radiation field mainly by photon-photon pair production and inverse Compton scattering:

$$\gamma + \gamma_{bb} \to e^+ + e^-, \quad e + \gamma_{bb} \to e' + \gamma' . \tag{3}$$

When cascade photons reach energies too low for pair production on the blackbody photons, the cascade development is slowed and further development occurs in the gas by ordinary pair production from γ scattering with nuclei of charge Z, but with electrons still losing energy mainly by inverse Compton scattering in the blackbody radiation

$$\gamma + Z \to Z + e^+ + e^-, \quad e + \gamma_{bb} \to e' + \gamma' . \tag{4}$$

At high energies, the cascade develops entirely on the blackbody photons by photon-photon pair production and inverse Compton scattering. The characteristic interaction rates for these processes are much higher than the expansion rate and thus one can assume the cascade spectrum is formed instantly. This spectrum is referred to as the ''zero-generation'' spectrum.

Since $\gamma - \gamma$ elastic scattering is the dominant process in a radiation-dominated plasma for photons just below the pair-production threshold [10], a primary photon or lepton triggers a cascade which develops until the photon energies have fallen below the maximum energy $E_{\max}(T) \simeq m_e^2/(22\,T)$ corresponding to the energy for which the mean-free paths against $\gamma - \gamma$ scattering and $\gamma - \gamma$ pair production are equal [11].

When $E_{\max}$ is less than the threshold for photodisintegration of ^{4}He nuclei (approximately 20 MeV), cascade nucleosynthesis is inefficient. This condition restricts the epoch of cascade nucleosynthesis to $T \lesssim 0.5$ keV. At temperatures between about 10^{-4} and 0.5 keV, where the subsequent cascade development is via ordinary pair production and inverse Compton scattering, the photons survive for a time determined by either the energy-loss rate for compton scattering or the interaction rate for ordinary pair production in the gas. During this time they can produce light nuclei by photodisintegration. The electrons and positrons give rise to first-generation photons as a result of inverse Compton scattering. These first-generation photons then produce the secondgeneration photons and so on. Each generation of photons is shifted to low energies because the inverse Compton scattering is in the Thompson regime and only one to two generations of photons are sufficiently energetic to induce cascade nucleosynthesis.

At $T \lesssim 10^{-4}$ keV, when interaction times become larger than the Hubble time, only the zero-generation photons are produced and the effectiveness of the cascade nucleosynthesis diminishes as T drops.

A detailed numerical cascade calculation including all the effects has been recently performed by Protheroe *et al.* [12] to find the number of deuterium and ^{3}He nuclei produced by the cascade initiated at a redshift z. It was shown that because of the almost instant formation of the zero-generation spectrum, the exact shape of the γ-ray or electron injection spectrum is of no consequence for further cascade development, and only the total amount of injected energy is relevant. The epoch of cascade nucleosynthesis is limited by $T_{\rm max}$ and $T_{\rm min}$. The maximum temperature is determined by the condition that the maximum energy of photons in the cascade spectrum, $E_{\rm max}$, must be larger than the threshold for D or ^{3}He production on ^{4}He, $E_{\rm max} \simeq 20$ MeV. This condition results in $T_{\rm max} \simeq 0.57$ keV. The effectiveness of D and ^{3}He production decreases as T decreases. The reason for that is not the decrease of the density of ^{4}He, but rather the decrease in number of low-energy photons in the cascade. In fact, the lower the temperature, the higher the photon energies in the cascade, and therefore a smaller fraction of the photons in cascade nucleosynthesis [12].

At $T \lesssim T_{\rm min} \simeq 2.4\times 10^{-4}$ keV the cascade photons can be directly observed and the upper limit for the isotropic γ-ray flux at $10-200$ MeV is more restrictive for the cascade production than nucleosynthesis. Therefore, the most effective epoch for cascade nucleosynthesis corresponds to temperatures in the range $10^{-4}-0.5$ keV [12].

It was also shown that the role of $\gamma\gamma \to \gamma\gamma$ scattering is important for epochs with temperatures $T \gtrsim 1.2\times 10^{-2}$ keV when the scattered photons do not interact again with the target photons, but are just redistributed over the spectrum producing a small bump before the cut-off energy $E_{\rm max}$.

IV. APPLICATION AND DISCUSSION

From the detailed analysis performed in Ref. [12], the number of ^{3}He and D nuclei produced by a single cascade of total energy E_0 is $\mathcal{N}(^3{\rm He}) \sim (0.1-1)E_0$ GeV^{-1} and $\mathcal{N}({\rm D}) \sim (1-6)\times 10^{-3} E_0$ GeV^{-1}.

Let us imagine that eternal self-annihilation of light photinos into lepton pairs are the only source of energetic cascades. In such a case, the total fraction $(^3{\rm He+D})/{\rm H}$ at the present epoch is

$$
\begin{aligned}
\frac{^3{\rm He}+{\rm D}}{\rm H} &= f_c\left(\frac{m_{\tilde\gamma}}{\rm GeV}\right)\int dt\,\frac{\dot n^c_{\tilde\gamma}}{n_{\rm H}}[\mathcal{N}(^3{\rm He},t)+\mathcal{N}({\rm D},t)] \\
&= -f_c\left(\frac{m_{\tilde\gamma}}{\rm GeV}\right)\int \frac{dt}{n_{\rm H}}[\mathcal{N}(^3{\rm He},t)+\mathcal{N}({\rm D},t)] \\
&\quad \times\langle\sigma_{\rm SA}|v|\rangle Y_\infty^2 s^2, \qquad (5)
\end{aligned}
$$

where f_c is the fraction of mass $m_{\tilde\gamma}$ transferred to cascade energy, $\mathcal{N}(^3{\rm He},t)$ and $\mathcal{N}({\rm D},t)$ are the number of ^{3}He and D nuclei produced at the time t per GeV by the total electromagnetic cascade [12], $n_{\rm H}$ is the hydrogen number density per comoving volume and $\dot n^c_{\tilde\gamma}$ represents the time derivative of the photino number density per comoving volume.

Photino self-annihilations in which the final state is a lepton-antilepton pair involve the t-channel exchange of a virtual slepton between the photinos, producing the final fermion-antifermion pair. In the low-energy limit the mass $m_{\tilde l}$ of the slepton is much greater than the total energy exchanged in the process and the photino-photino-fermion-antifermion operator appears in the low-energy theory with a coefficient proportional to $e_i^2/m^2_{\tilde l}$, with e_i the charge of the final-state fermion. Also, as there are two identical fermions in the initial state, the annihilation proceeds as a p wave, which introduces a factor v^2 in the low-energy cross section [13]. The resultant low-energy photino self-annihilation cross section is [14]

$$
\begin{aligned}
\langle\sigma_{\rm SA}|v|\rangle &= 8\pi\alpha^2_{\rm em}\sum_i e_i^4\frac{m^2_{\tilde\gamma}}{m^4_{\tilde l}}\frac{v^2}{3} \simeq 2.3\times 10^{-11}\left(\frac{T}{m_{\tilde\gamma}}\right) \\
&\quad\times\left(\frac{m_{\tilde\gamma}}{\rm GeV}\right)^2\left(\frac{m_{\tilde l}}{100\ {\rm GeV}}\right)^{-4}{\rm mb}, \qquad (6)
\end{aligned}
$$

where we have used for the relative velocity $v^2 = 6/x$ and we have summed over e and μ assuming a common slepton mass $m_{\tilde l}$ for selectron and smuon scalar fields.

Recalling that for a radiation-dominated Universe $t = 0.301 g_*^{-1/2}(M_{\rm Pl}/T^2)$, assuming the baryonic mass is 77% hydrogen by mass, and making use of Eq. (2), we have numerically evaluated the time integral of the right-hand side of Eq. (5) and obtained the present total fraction $(^3{\rm He+D})/{\rm H}$ generated by the electromagnetic showers induced by eternal photino self-annihilations into lepton-antilepton pairs:

$$
\frac{^3{\rm He}+{\rm D}}{\rm H} \simeq 9.4\times 10^{-12}\left(\frac{f_c}{0.77\Omega_{\rm B}}\right)(\Omega_{\tilde\gamma}h)^2\left(\frac{m_{\tilde l}}{100\ {\rm GeV}}\right)^{-4}. \qquad (7)
$$

The upper limit inferred from measurements of ^{3}He in meteorites and the solar wind (which requires making assumptions about stellar processing and galactic chemical evolution) is $(^3{\rm He+D})/{\rm H} \lesssim 1.1\times 10^{-4}$ [15]. This bound translates from Eq. (7) into a lower limit on the slepton mass[2]

$$
m_{\tilde l} \gtrsim 1.7\left(\frac{f_c}{0.77\Omega_{\rm B}}\right)^{1/4}(\Omega_{\tilde\gamma}h)^{1/2}\ {\rm GeV}. \qquad (8)
$$

We notice that this lower bound on the slepton mass is *independent* of the photino mass and thus holds for the entire range of cosmological interest of the photino mass

[2]Note, however, that very recent measurements of D/H are closer to $(2-5)\times 10^{-4}$ [16] and, if such higher values for $(^3{\rm He+D})/{\rm H}$ were adopted, the upper limits we derive would be correspondingly higher.

$1.6 \lesssim r \lesssim 2$ [3]. Since the lower limit on $m_{\tilde{l}}$ given in Eq. (8) is weaker than any other present experimental lower bound on slepton masses, we may conclude that eternal self-annihilations of light photinos do not jeopardize the successful predictions of primordial nucleosynthesis.

Now let us generalize our analysis to the case in which the relic abundance of a generic low-mass dark-matter candidate X is determined by some processes other than its self-annihilation cross section into leptons, which we parametrize as $\langle \sigma_{\rm SA} |v| \rangle = \sigma_0 (T/m_X)^n$, where we have assumed that either s-wave ($n=0$) or p-wave ($n=1$) annihilations dominate;[3] it is straightforward to generalize to the case where both contribute. For instance, one can envisage the situation in which the dark-matter component X is slightly degenerate in mass with a particle X' and that interconversion processes $X \leftrightarrow X'$ keep X in equilibrium even after self-annihilations have frozen out, regardless of the strength of the interactions governing the X' abundance. This case is rather different from the one analyzed in Ref. [5] where it was assumed that the same cross section determines both the self-annihilation rate and the relic abundance of the X particles.

By making use of Eqs. (2) and (5), we have evaluated numerically the present fraction of (^{3}He+D)/H generated by electromagnetic cascades induced by eternal self-annihilations of X's into leptons. Adopting the upper limit (^{3}He+D)/H$\lesssim 1.1 \times 10^{-4}$ [15] we find

$$\left(\frac{\sigma_0}{10^{-38}\mathrm{cm}^{-2}}\right)\left(\frac{m_X}{\mathrm{GeV}}\right)^{-1} \lesssim 2.7\times 10^{7}\left(\frac{f_c}{0.77\Omega_{\rm B}}\right)^{-1}(\Omega_X h)^{-2} \quad \text{for } n=0,$$

$$\left(\frac{\sigma_0}{10^{-38}\mathrm{cm}^{-2}}\right)\left(\frac{m_X}{\mathrm{GeV}}\right)^{-2} \lesssim 6.6\times 10^{12}\left(\frac{f_c}{0.77\Omega_{\rm B}}\right)^{-1}(\Omega_X h)^{-2} \quad \text{for } n=1. \tag{9}$$

Once the model-dependent cross section σ_0 is known, one can apply the above limits to find the appropriate bounds on the parameter space of the model.

So far we have been assuming that the generic dark-matter component X eternally self-annihilates only into lepton-antilepton pairs. When hadronic channels are open, D is produced by hadronic showers and this requires reconsideration of the constraint derived above. In fact, even if light photinos self-annihilate only into lepton pairs, the resulting electromagnetic cascades will be effectively hadronic if $E_\gamma \varepsilon_\gamma \gtrsim$ GeV^2, where E_γ is the typical energy of a cascade photon and ε_γ is the energy of a blackbody photon. Furthermore, there is always about a 1% probability for the virtual decay photon to convert into a quark-antiquark pair over threshold. Hence, hadronic showers will be generated if $m_X \gtrsim 1$ GeV even if X-particles do not have specific self-annihilation channels into quark pairs (see Reno and Seckel [5]). Let us analyze the possible different channels in more details in the case of a light photino.

For $m_{\tilde{\gamma}} \gtrsim 1$ GeV pions may be produced copiusly since we are above (multi) pion threshold. Pions generated at temperatures $T \ll 1$ MeV decay without thermalizing with the surrounding gas and produce energetic photons and electrons [5]. This case, therefore, reduces to what was studied in the previous sections. However, for temperatures greater than about 1 MeV the strong interaction rate is competitive, or at least non-negligible, compared to the pion decay rate. Therefore, long-lived pions may change the neutron-to-proton ratio through processes such as $\pi^- p \to n \pi^0$ at the time (or just after) the ratio is being fixed by weak interactions, resulting in a shift of the primordial helium abundance. We can make an estimate of this shift in the following way. Making use of Eq. (1), it is easy to show that the number of annihilations of photinos per baryon is

$$\frac{\Delta Y_{\tilde{\gamma}}}{Y_B} = 2.3\times 10^{-9} g_*^{-1/2} g_{*S} \frac{T_1^2 - T_2^2}{m_{\tilde{\gamma}} Y_B} \ \mathrm{GeV}^{-1}. \tag{10}$$

At the time the neutron-to-proton ratio is being fixed by the weak interactions the temperature is about 0.5 MeV. Taking $T_1 - T_2 \sim T_1 \sim 0.5$ MeV, we find $\Delta Y_{\tilde{\gamma}}/Y_B \sim 10^{-7}$. So, following Seckel and Reno [5], we expect pions created by light-photino annihilations to have a very small effect upon the helium abundance. Since the deuterium abundance is set even later, pion interactions also have only a small effect on the deuterium abundance.

If photinos annihilate into protons and neutrons, the effects on the element abundances are more dramatic. As discussed in detail by Dimopoulos *et al.* [17], the main effect is the destruction of the ambient ^{4}He and the creation of D (which is photodestroyed by electromagnetic showers), ^{3}He, ^{6}Li, and ^{7}Li. Even though the final light-element abundances may be found only by an extensive numerical analysis involving both hadronic and electromagnetic cascades, the constraints on the primordial abundance of ^{6}Li and ^{7}Li are so stringent [9] one might expect that the energy released into nucleons by self-annihilation of photinos must be very small. Small numbers of nucleons may result either if the photino is sufficiently light so that emission of quark pairs is phase-space suppressed, or if the self-annihilation cross section is small, which may happen if the virtual particles mediating the process are very heavy.

In conclusion, we have investigated the effect on big-bang nucleosynthesis predictions of a very light photino predicted in low-energy supersymmetry models in which dimension-3 supersymmetry breaking operators are suppressed [2]. We have found that eternal photino self-annihilations into fermion pairs that take place around 10^5 to 10^6 sec after the big-bang cannot significantly alter the primordial abundances of light elements when taking into account the present experimental limits on the slepton masses.

[3]This may happen, for instance, when explicit sfermion mixing [18] or CP-violating phases [19] in low-energy supersymmetric models greatly enhance the self-annihilation cross-section of the lightest supersymmetric particles into fermions removing the p-wave suppression.

We have also generalized our study to the case of a generic dark-matter component X whose relic abundance is not determined by the self-annihilation cross section into leptons, but by some other processes, e.g., co-annihilations with a slightly heavier particle X'. Limits on the self-annihilation cross section and mass of the X particle have been given by requiring that electromagnetic cascades generated by the eternal self-annihilations of the X particle do not jeopardize the predictions of the big-bang nucleosynthesis.

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