TeV $\gamma$ rays from photodisintegration and daughter deexcitation of cosmic-ray nuclei

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Introduction. In the field of TeV $\gamma$-ray-astronomy, new instruments are discovering new sources at a rapid rate, both within our Galaxy and outside the Galaxy\cite{1}. Not surprisingly, one of the brightest TeV $\gamma$-ray sources is the one discovered first, the Crab pulsar wind nebula. The integral $\gamma$-ray flux obtained from the Crab by the Whipple Collaboration is now the standard TeV flux unit: $F_{\text{Crab}}(\gamma_{\text{LAB}} > 0.35 \text{TeV}) = 10^{-10} / \text{cm}^2 / \text{s}$\cite{2}. The spectral index of the Crab’s integrated flux is measured to be $-1.5$, so $F_{\text{Crab}}$ falls by a factor of 30 for each decade of increase in $\gamma_{\text{LAB}}$. Newly commissioned atmospheric Cerenkov telescopes (CANGAROO, HESS, MAGIC and VERITAS) will reach a sensitivity 100 times below $F_{\text{Crab}}$.

Two well-known mechanisms for generating TeV $\gamma$-rays in astrophysical sources\cite{3} are the purely electromagnetic (EM) synchrotron emission and inverse Compton scattering, and the hadronic (PION) one in which $\gamma$-rays originate from $p\gamma$ production and decay. There has been considerable debate over which of these two mechanisms is dominant. The mechanism for the PION mode may be either $pp$ or $p\gamma$ collisions. In this Letter, we highlight a third dynamic which leads to TeV $\gamma$-ray emission via the photo-disintegration of highly boosted nuclei followed by daughter de-excitation. Starburst regions such as Cygnus OB2 appear to be promising sites for TeV $\gamma$-ray emission via this mechanism. A unique feature of the $A^*$ process is a sharp flattening of the energy spectrum below $\sim 10 \text{TeV}/(T/\text{eV})$ for $\gamma$-ray emission from a thermal region of temperature $T$. We also check that a diffuse $\gamma$-ray component resulting from the interaction of a possible extreme-energy cosmic-ray nuclei with background radiation is well below the observed EGRET data. The $A^*$ mechanism described herein offers an important contribution to $\gamma$-ray astronomy in the era of intense observational activity.

astrophysical context was first calculated by Stecker in a seminal paper\cite{4}, almost forty years ago. To our knowledge, the possibly important role of nuclear de-excitation in the astrophysical context was first appreciated and proposed by Moskalenko and collaborators\cite{5} more than a decade ago. Since then, the $A^*$-process has been overlooked by the $\gamma$-ray community. Now that data is becoming available which can validate or invalidate the process, it is timely to revive and further develop the $A^*$-process. We do so, by providing calculational details, and by identifying the astrophysical context where this process might dominate over the EM and PION modes for production of $\gamma$-rays. A detailed discussion on the main features of the EM and PION mechanisms is presented in a longer accompanying paper\cite{6}.

By far the largest contribution to the photo-excitation cross-section comes from the Giant Dipole Resonance (GDR) at $\epsilon_{\text{GDR}} \sim 10 \text{ MeV} - 30 \text{ MeV}$ in the nuclear rest frame\cite{7}. The ambient photon energy required to excite the GDR is therefore $\epsilon = \epsilon_{\text{GDR}} / \Gamma_A$. The GDR decays by the statistical emission of a single nucleon, leaving an excited daughter nucleus $(A-1)^*$. The probability for de-excitation by emitting one or more photons of energies $\epsilon_{\text{GDR}} \sim 1 - 5 \text{ MeV}$ in the nuclear rest frame. The lab-frame energy of the $\gamma$-ray is then $\epsilon_{\gamma}^{\text{LAB}} = \Gamma_A \epsilon_{\gamma}^{\text{GDR}}$.

As just outlined, the boost in the nuclear energy from rest to $E_{\text{LAB}}^A = \Gamma_A m_N$ plays two roles. It promotes the thermal energies of the ambient photons to the tens of MeV $\gamma$-rays capable of exciting the GDR, and it promotes the de-excitation photons from a few MeV to much higher energies, potentially detectable in $\gamma$-ray telescopes. Eliminating $\Gamma_A$ above leads to the relation $\epsilon_{\gamma}^{\text{LAB}} \sim \epsilon_{\gamma}^{\text{GDR}} \epsilon_{\gamma}^{\text{dm}}$. The "\sim" indicates that the $\epsilon_{\gamma}^{\text{GDR}}$ and $\epsilon_{\gamma}^{\text{dm}}$ energies are not sharp, but rather distributed over their resonant shapes ("Lorentzian" or
“Breit-Wigner”). Each of the two spectra have a width to mass ratio less than unity. It is sufficient for our purposes to treat each spectrum in the narrow-width-approximation (NWA). Well-defined values for $\varepsilon^{\text{GDR}}$ and $\epsilon^{\text{dnn}}$ then result. We may summarize up to this point by saying that the $A^*$ process produces $\gamma$-ray’s with energy $E^{\text{LAB}}_{\gamma} = \varepsilon^{\text{GDR}}/\epsilon^{\text{dnn}} \sim 20 \text{TeV}/(\text{T/ev})$ if there exists an accelerated nuclear flux with boost $\Gamma_A \sim \varepsilon^{GDR}/\epsilon \sim 7 \times 10^6 (\text{T/ev})$ or equivalent energy $E^{\text{LAB}}_{\gamma} \sim 7 \text{PeV}/(\text{T/ev})$ per nucleon [6, 7].

Here and throughout we assume a Bose-Einstein (BE) distribution with temperature $T$ for the ambient photons; in the mean, $\langle \epsilon \rangle \sim 3.3 T$. As an example, $\gamma$-ray’s with energy $E^{\text{LAB}}_{\gamma} \sim 20 \text{TeV}$ are generated in this $A^*$-process if an ambient photon temperature of an eV and a boost factor of $7 \times 10^6$ are present. Thus, we have arrived at an astrophysical environment where the $A^*$-process may dominate production of TeV $\gamma$-ray’s: the region must contain far-UV photons (commonly defined as 1-20 eV) from the Lyman-$\alpha$ emission of young, massive, hot stars such as O and B stars which have surface temperatures of $T_{\odot} \sim 40,000 \text{K} (3.4 \text{eV}$ and $\langle \epsilon \rangle \sim 10 \text{ eV})$ and $18,000 \text{K} (1.5 \text{ eV}$ and $\langle \epsilon \rangle \sim 5 \text{ eV})$, respectively; and the region must contain shocks, giant winds, or other mechanisms which accelerate nuclei to excess of a PeV per nucleon. Violent starburst regions, such as the one in the direction of Cygnus, are splendid examples of regions which contain OB stars, shocks and giant stellar winds.

With $\epsilon^{GDR}$ and $\epsilon^{dnn}$ fixed at their central values, the $\gamma$-ray spectrum and $\text{d}n(\epsilon^{LAB})/\text{d}\epsilon^{LAB}$ is given by a simple Jacobian times the BE distribution with argument $\epsilon$ set to $\epsilon^{GDR}/\epsilon^{dnn}$. i.e.,

$$\frac{\text{d}n(\epsilon^{LAB})}{\text{d}\epsilon^{LAB}} \propto (\epsilon^{LAB})^{-4} \left( \epsilon^{GDR}/\epsilon^{dnn} - \epsilon^{LAB} \right)^{-1} \cdot (1)$$

Three regions in $\epsilon^{LAB}$ emerge. For $\epsilon^{LAB} < \epsilon^{GDR}/\epsilon^{dnn}$, the $\gamma$-ray spectrum is exponentially suppressed [11]. For $\epsilon^{LAB} \gg \epsilon^{GDR}/\epsilon^{dnn}$, there is power-law suppression. The $\gamma$-ray spectrum peaks near $\epsilon^{LAB} \sim \epsilon^{GDR}/\epsilon^{dnn}/T$. In particular, one notes that the suppression of the high-energy Wien end of the thermal photon spectrum has led to a similar suppression of the photon spectrum below $\epsilon^{LAB} \sim 20 \text{ TeV}/(\text{T/ev})$. This lower-energy suppression presents a robust prediction of the $A^*$-process. Moreover, it contrasts greatly with the PION and EM processes, and so provides a unique signature. The $A^*$ process predicts “orphan” sources, with suppression of associated GeV $\gamma$-ray’s or MeV X-rays (although the photo-dissociated neutrons may $\beta$-decay to neutrinos [10]). Observationally, orphan $\gamma$-ray sources have been identified [12, 13] and some orphans are known to be near OB starburst regions [14].

Comparing to the PION $p\gamma$ process, the $A^*$ process has a much lower ambient photon energy threshold; in the nuclear rest frame, the photon threshold is $\varepsilon^{GDR} \sim 10 \text{ MeV}$ for the $A^*$ process, but it is an order of magnitude larger at $m_{\pi}$ for the PION $p\gamma$ process [16]. This means that for fixed ambient temperature, $\Gamma_A$ need be an order of magnitude larger for the PION $p\gamma$ process, and the resulting $\gamma$-ray energies, proportional to $\Gamma_A^2$, are two orders of magnitude larger. The EM and PION $pp$ processes contrast with the $A^*$-process in that there is either no threshold (EM) or very small threshold $\mathcal{O}(2m_{\pi})$ in the lab (PION $pp$), and the resulting $\gamma$-ray spectrum rises monotonically with decreasing $\epsilon^{LAB}$.

The $A^*$-Rate. We now derive the rate for $\gamma$-ray production in the $A^*$-process. The cross-section is dominated by the GDR dipole form [10], which in the NWA is

$$\sigma_A(\epsilon) \rightarrow \frac{\pi}{2} \sigma^{\text{GDR}} \Gamma^{\text{GDR}} \delta(\epsilon^{\text{GDR}} - \Gamma_A \epsilon) \cdot (2)$$

where $\Gamma^{\text{GDR}}$ and $\sigma^{\text{GDR}}$ are the GDR width and cross-section at maximum. Fitted numerical formulas are $\sigma^{\text{GDR}} = 1.45 A \times 10^{-2} \text{cm}^2$, $\Gamma^{\text{GDR}} = 8 \text{ MeV}$, and $\epsilon^{\text{GDR}} \sim 42.65 A^{0.41}$ MeV for $A > 4$ and $0.925 A^{2.43}$ for $A \leq 4$ [17]. The 8 MeV width implies a very short de-excitation distance of $\sim 25 \Gamma_A \text{fm}$.

The general formula for the inverse photo-disintegration mean-free-path (mfp) [18] for a highly relativistic nucleus with energy $\Gamma_A/m_A$ propagating through an isotropic photon background with energy $\epsilon$ and spectrum $dn(\epsilon)/d\epsilon$ is [5, 7]

$$(\lambda_A)^{-1} \rightarrow \frac{\pi \sigma^{\text{GDR}} \Gamma^{\text{GDR}}}{4 \Gamma_A^2} \int_{\epsilon^{\text{GDR}}}^{\epsilon} \frac{d\epsilon}{\epsilon^2} \frac{dn(\epsilon)}{d\epsilon} \cdot (3)$$

For a nucleus passing through a region of thermal photons, integration over the BE photon distribution gives

$$(\lambda_A^{\text{BE}})^{-1} \approx \frac{\pi \sigma^{\text{GDR}} \Gamma^{\text{GDR}}}{\epsilon^{\text{GDR}}} \int_{\epsilon^{\text{dnn}}}^{\epsilon_{\gamma}} d\epsilon \ln \left( 1 - \frac{e^{-w}}{w} \right) \cdot (4)$$

where we have defined a dimensionless scaling variable $w = \epsilon^{\text{GDR}}/2 \Gamma_A T$. From the prefactor, we learn that the peak of $(\lambda_A^{\text{BE}})^{-1}$ scales in $A$ as $\sigma^{\text{GDR}}/\epsilon^{\text{GDR}} \sim A^{1.21}$, and the value of $\Gamma_A$ at the peak scales as $\epsilon^{\text{GDR}} \sim A^{-0.21}$.

The scaling function $f(w) = w^2 \ln (1 - e^{-w})$ is shown in Fig. [1]. Approximations to the $|\ln w|$ term yield $e^{-w}$ for $w > 2$, and $|\ln w|$ for $w < 1$. Thus, the exponential suppression of the process appears again for large for $w > 2$, i.e., for $\epsilon^{LAB} < \epsilon^{GDR}/\epsilon^{dnn}/4T$, and the small $w$ region presents a mfp that scales as $w^2 |\ln w|$. The peak region provides the smallest inverse mfp, and so this region dominates the $A^*$-process. In the peak region, $w$ is of order one, which implies that $\Gamma_A T \sim \epsilon^{\text{GDR}}$. When this latter relation between the nuclei boost and the ambient photon temperature is met, then the photo-disintegration rate is optimized.

The area under the peak region in Fig. [1] is of order one, which leads to a simple and reasonable estimate of
the inverse mfp given in Eq. (4), for any nucleus boosted near $\Gamma_A \sim \epsilon^\text{GDR}/T$. The useful and eminently sensible estimate is $(\lambda^\text{BE}_A)^{-1} \sim \sigma^\text{GDR} \lambda^\text{BE}_A$; here, we have input the typical value $\Gamma^\text{GDR}/\epsilon^\text{GDR} \sim 1/3$. Putting in numbers, this estimate yields

$$\lambda^\text{BE}_A \sim \frac{5 \times 10^{13} \text{ cm}}{A \ (\text{T/eV})^3} \sim \frac{3 \text{ AU}}{A \ (\text{T/eV})^3} \tag{5}$$

for a nucleus with energy in the peak regime around $E_A \sim 10 A/(\text{T/eV}) \text{ PeV}$. Here we have used $n^\text{BE}_\gamma = 2 \zeta(3) T^3/\pi^2 \simeq T^3/4$.

An important question is how many photo-disintegration steps in the nuclear chain $A \to (A-1)^* \to (A-2)^*$ etc. may be expected. Each step will produce an excited daughter which may de-excite via emission of $\gamma$-ray’s, each having a typical lab-frame energy $\epsilon^\text{LAB}_\gamma \sim \epsilon^\text{GDR}_\gamma \epsilon^\text{dSN}/3T$. Clearly, the number of steps depends on the mfp of each excited daughter nucleus, and on the diffusion time of the nuclei in the thermal region. The probability for $n$-cascades, producing $n$ excited daughter nuclei and $n$ single nucleons, may be written in symmetric ways as

$$P_n = \prod_{j=1}^n \int_0^{2x_{j+1}/\lambda_j} \frac{dx_j}{\lambda_j} e^{-\epsilon_{x_j-x_{j-1}}} = \prod_{j=1}^n \int_0^{x_{j+1}/\lambda_j} dx_j e^{-x_j \delta_j},$$

where $x_0 = 0$, $x_{n+1}$ equals the diffusion length $D$ of nuclei in the $A^*$-region, the ordered $x_1 \leq x_2 \ldots \leq x_n$ denote the spatial positions of successive photo-disintegrations to $(A-1)^*, \ldots, (A-n)^*$, $\delta_j = \lambda_j^{-1} - \lambda_{j+1}^{-1}$, and $\lambda_{n+1} = \infty$. The exponentials in (6) are probabilities that the various daughters not interact (i.e., survive) from the point of creation to the point of photo-disintegration. A simple result obtains when the mfp’s are long on the diffusion scale $D$ of the $A^*$-region. In this case the nested integrals collapse to $P_n \approx (n!)^{-1} \prod_{j=1}^n D/\lambda_j$. This is just the product of the independent probabilities for each excited daughter to be produced, times the factor $1/n!$ that divides out all but the one correct time ordering of the $n$ photo-disintegrations. One sees that, since $D/\lambda_j \ll 1$ by assumption, disintegration to more than the first excited daughter is unlikely. Such is the case if photo-disintegration of nuclei occurs in the Cygnus OB starburst region, as we now demonstrate.

**Relevance to Starburst Regions.** The aforementioned possibility of generating Galactic TeV $\gamma$-ray’s from accelerated nuclei scattering on starlight in starburst regions is of considerable astrophysical interest. The energy of the photons ($\sim 3T^*$ in the mean) is maintained as the photons disperse from the stars, but the photon density decreases by the ratio of the surface of stellar emission ($N_* \times 4\pi R^2_\ast$) to the loss surface of the starburst region ($4\pi R^2_{\text{SB}}$). Taking the giant star radius to be $R_* \sim 10 R_\odot$, the radius of the starburst region to be $R_{\text{SB}} \sim 10$ pc, and the stellar count to be $N_* \sim 2600$, the photon density is diluted by $\sim 10^{12}$. According to Eqs. (6) and (5), $(\lambda^A)^{-1}$ is proportional to the photon density, and hence to this factor. Including this factor in Eq. (5), one gets the estimate $\lambda^A \sim (56/A) (1.5 \text{ eV}/T_\ast) \times 10^{23} \text{ cm}$ for the nuclear mfp in the starburst region. The hot stars are dominantly (95%) B type, with $T_\ast \sim 1.5 \text{ eV}$.

We note here that the diffusion time for a nucleus in the Cygnus OB starburst region is calculated to be $\sim 10^4 \text{ yr} \times 10^{22} \text{ cm}$, an order of magnitude shorter than the nuclear mfp. Accordingly, one expects only a few percent of the nuclei to photo-disintegrate in the Cygnus OB region, with multiple disintegrations being rather rare.

Some early statistical-model calculations for the production of $\gamma$-rays through the decay of the GDR gave a mean photon multiplicity between 0.5 and 2 for the new nuclei in the chain reaction (see Ref. [7] for details). We will simplify the data by assuming that one photon is emitted per nuclear de-excitation. Then, the rate of photo-emission is just the rate of excited daughter production, which in turn is just the rate of photodisintegration. Allowing for the $1/r^2$ dilution of the $\gamma$-ray flux from the source region of volume $V_A$, at distance $d$, the observed integral $\gamma$-ray flux $F_\gamma$ at Earth with $\epsilon^\text{LAB}_\gamma$ above a few TeV $= (\text{eV}/T)$ is then

$$F_\gamma = \frac{V_A}{4\pi d^2} \frac{1}{\lambda^A} F_A,$$

where the flux $F_A(= c n_A)$ is integral over the energy-decay of the peak region, $A (\text{eV/T}) \text{ PeV} < E_A < 10 A (\text{eV}/T)$. It is commonly assumed that this nuclear flux results from continuous trapping of the diffuse cosmic ray flux by diffusion in a milligauss magnetic field.

Putting into Eqs. (5) and (7) the parameters for the Cygnus OB2 region, i.e., $\lambda_A = (56/A) \times 10^{23} \text{ cm}$, $R_{\text{SB}} =...
10 pc, and $d = 1.7$ kpc, one obtains $F_\gamma = 2 \times 10^{-10} A F_A$. Thus, an accumulated nuclear flux within Cygnus OB2 of $f/\text{cm}^2/\text{s}$ in the PeV region gives $A \geq \text{TeV} \gamma$-ray flux at Earth of $\sim 10 A f \times F_{\text{Crab}} (\epsilon^{\text{LAB}}_\gamma \geq 1 \text{ TeV})$. For example, an iron flux $F_{\gamma 56} = 6 \times 10^{-5} \text{cm}^2/\text{s}$ above a TeV leads to a $\gamma$-ray flux at Earth of about 3% of the Crab, just at the level measured by the HEGRA experiment for $\gamma$-ray’s from the Cygnus OB2 direction [12]. In our accompanying paper [7] we present a more detailed calculation of the $A^*$-process for this starburst region, and show that $F_{\gamma 56} \sim 10^{-4}/\text{cm}^2/\text{s}$ is a credible 1% of the kinetic energy budget of Cygnus OB2.

The $\gamma$-ray energy spectrum remains to be discussed. In the rest frame of the excited nucleus, the photon is emitted isotropically and nearly monochromatically. Therefore, in the lab frame, the $\gamma$-ray-spectrum is nearly flat between 0 and $2 \Gamma A \epsilon^{\text{xn}}_\gamma$ on a linear scale. The power spectrum and integrated power rise as $\epsilon^{\text{LAB}}_\gamma$ and $(\epsilon^{\text{LAB}}_\gamma)^2$, respectively, peaking near $\epsilon^{\text{LAB}}_\gamma \sim \epsilon^{\text{GDR}}_\gamma \epsilon^{\text{xn}}_\gamma/4 \Gamma \sim 10 \text{ TeV (eV/T)}$, as explained in the paragraph below Eq. (4).

In detail, the photon spectrum is obtained by replacing the approximate Eq. (7) with an integral over $(dF_A/d\epsilon^{\text{LAB}}_\gamma) \times (\lambda^{3A})^{-1}$, with measure $d\epsilon^{\text{LAB}}_\gamma \cos \theta_\gamma \delta(\epsilon^{\text{LAB}}_\gamma - \Gamma A \epsilon^{\text{xn}}_\gamma (1 + \cos \theta_\gamma))$, where $\theta_\gamma$ is the angle in the nucleus rest frame between the isotropically-emitted photon and the boost direction. After assuming a power-law (with spectral index $\alpha$) for the nuclear spectrum, and integrating over $d\cos \theta_\gamma$, one arrives at

$$\frac{(\epsilon^{\text{LAB}}_\gamma)^2 dF_\gamma}{d\epsilon^{\text{LAB}}_\gamma} = \frac{n^{\text{Th}} V_{A^*} \alpha_{\text{GDR}} \Gamma_{\text{GDR}} \epsilon^{\text{xn}}_\gamma m_N}{4 \pi d^2} \times \left( \epsilon^{\text{GDR}}_\gamma \right)^{\alpha - 2} \left( \frac{2E_N}{m_N} \right)^{\alpha} \frac{dE_N}{dE_N} x^2 F(x_e),$$

where $n^{\text{Th}}$ is the true density of the thermal photons after spatial dilution, $E_0$ is any reference energy, $x_e = \epsilon^{\text{LAB}}_\gamma T/\epsilon^{\text{xn}}_\gamma \epsilon^{\text{GDR}}_\gamma$ is the dimensionless energy-scaling variable, and

$$F(x_e) = \int \frac{d\omega}{\omega} \omega^\alpha f(\omega) = \int_0^{1/x_e} d\omega \omega^{1+\alpha} \left| \ln(1 - e^{-\omega}) \right|$$

is the scaling function. Shown in Fig. 1 is $x_e^2 F(x_e)$ [15]. The $\gamma$-ray energy is $\epsilon^{\text{LAB}}_\gamma = 10 x_e \text{TeV}$ $\times$ $(\epsilon^{\text{xn}}_\gamma/\text{MeV}) (\epsilon^{\text{GDR}}_\gamma/10 \text{MeV}) (\text{T/eV})^{-1}$. We note that the prefactor of $x_e^2 F(x_e)$ in Eq. (5) scales in $A$ for fixed energy per nucleon as $\sigma^{\text{GDR}}_\gamma (\epsilon^{\text{GDR}}_\gamma)^{1+0.21(\alpha-1)}$, while the position of the peak in the $x_e^2 F(x_e)$ spectrum at $x_e \sim 0.25$ scales mildly as $x_e \propto \epsilon^{\text{GDR}}_\gamma A^{1+0.21}$. Well beyond the peak, the $x_e^2 F(x_e)$ spectrum falls as $x_e^{-2} \ln x_e$; equivalently, the $\gamma$-ray spectrum falls as $\sim (\epsilon^{\text{LAB}}_\gamma)^{-4}$.

The Cosmogenic $A^*$-Process. One well-known application of nuclear photo-disintegration occurs in the calculation of the propagation of extreme-energy cosmic nuclei in the cosmic microwave background (CMB) and cosmic infrared background (CIRB). The CMB contribution is thermal, with temperature $T = 2.3 \times 10^{-6} \text{ eV}$, and Eq. (3) readily yields a mfp of $A^{-1}$ Mpc for nuclei with energies near $E_A \sim 4 A \times 10^{19} \text{ eV}$. The CIRB, now fairly well known [19], is much smaller than the CMB. Its spectrum may be approximated by one or two power-laws, which allows an analytic integration of Eq. (3). The result is a mfp longer than that of the CMB result, but relevant to nuclei with energies down to $10^{16} \text{ eV}$.

Of interest in our work is the $\gamma$-ray flux produced when the photo-dissociated nuclear fragments produced on the CMB and CIRB de-excite. These $\gamma$-rays create chains of electromagnetic cascades on the CMB and CIRB, resulting in a transfer of the initial energy into the so-called EGRET region below 100 GeV, which is bounded by observation to not exceed $\omega_{\text{max}} \sim 2 \times 10^{-6} \text{ eV/cm}^3$ [20].

Fortunately, we can finesse the details of the calculation by arguing in analogy to work already done. Two recent papers [21, 22] have calculated the $\bar{\nu}_e$ flux resulting from photo-disintegration of cosmic Fe, followed by $\beta$-decay of the associated free neutrons. The photo-disintegration chain produces one $\beta$-decay neutrino with energy of order 0.5 MeV in the nuclear rest frame, for each neutron produced. Multiplying this result by 2 to include photo-disintegration to protons in addition to neutrons correctly weights the number of steps in the chain. Each step produces on average one photon with energy $\sim 3 \text{ MeV}$ in the nuclear rest frame. Comparing, about 12 times more energy is deposited into photons. Including the factor of 12 relating $\omega_{\text{max}}$ to $\omega_{\nu}$, we find from Fig. 3 in Ref. [21] that cosmogenic photo-disintegration/de-excitation energy is more than three orders of magnitude below the EGRET bound [23]. This result appears to be nearly invariant with respect to varying the maximum energy of the Fe injection spectrum (with a larger $E_{\text{max}}$, the additional energy goes into cosmogenic pion production). Thus, there is no constraint on a heavy nuclei cosmic-ray flux from the $A^*$ mechanism.

Conclusion. In final summary, we have presented an alternative mechanism ($A^*$) for generating TeV $\gamma$-ray’s in starburst regions. It has a unique orphan-like signature, with a flat spectrum below $\sim 20 \text{ TeV/(T/eV)}$ and a quasi power-law above. Since the EM and PION mechanisms also produce unique signatures, gamma-ray astrophysics benefits from a one-to-one correspondence between the source dynamic and the observable spectrum.

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