A possible explanation of the knee of cosmic light component spectrum from 100 TeV to 3 PeV*

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Abstract: The mixed Hydrogen and Helium (H + He) spectrum with a clear steepeening at $\sim 700$ TeV has been detected by ARGO-YBJ experiments. In this paper, we demonstrate that the observed H + He spectrum can be well reproduced with the model of cosmic rays escaping from the supernova remnants (SNRs) in our Galaxy. In this model, particles are accelerated in a SNR through a non-linear diffusive shock acceleration mechanism and three components of high energy light nuclei escaped from the SNR are considered. It should be noted that the proton spectrum observed by KASCADE can be also explained by this model given a higher acceleration efficiency.

Key words: Cosmic Rays, particle acceleration, supernova remnants, amplified magnetic field

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1 Introduction

Supernova remnants (SNRs) are generally believed to be the origin of Galactic cosmic rays (CRs)\textsuperscript{[1]}\textsuperscript{[5]}. The knee of $\sim 3 \times 10^{15}$ eV as a feature coinciding with the maximum energy of the light component of cosmic rays and the transition to a gradually heavier mass composition are mainly based on KASCADE results\textsuperscript{[16]}. Some recent data, however, appear to challenge this finding: the combined detection of showers with a wide field of view Cherenkov telescope (WFCT) and ARGO-YBJ find a flux reduction in the light component at $\sim 700$ TeV\textsuperscript{[7]}. A possible explanation of the knee of cosmic light component spectrum observed by KASCADE can be also explained by this model given a higher acceleration efficiency.

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2 The Review of NLDSA Model

In the NLDSA model given by Ref.\textsuperscript{[15]}, the shock is placed at a distance $x = 0$, so the upstream and downstream region correspond to $x < 0$ and $x > 0$, respectively. Physical quantities measured at upstream infinity, immediate upstream of the shock, and downstream are labeled by subscripts 0, 1, and 2, respectively. Generally, it is convenient to define two different compression ratios for subshock and total shock\textsuperscript{[19]}: $r_{\text{sub}} = \tilde{u}_1/\tilde{u}_2$ and $r_{\text{tot}} = \tilde{u}_0/\tilde{u}_2$, where $\tilde{u} = u + u_A$, $u$ is the bulk plasma velocity in the shock frame, and $u_A$ is the Alfvén velocity with respect to the background plasma.

In such a system, a diffusive-convection equation describing the transport of the $i$th particles in the shock

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frame can be expressed as [13]:
\[
\frac{\partial f}{\partial x} = \frac{p \partial u}{3 \partial x} \frac{\partial f}{\partial p} + \frac{\partial}{\partial x} \left[ \kappa_i(x,p) \frac{\partial f}{\partial x} \right] + Q_i(x,p),
\]
(1)
where subscript \( i \) represents the \( i \)th particles, \( f_i = f_i(x,p) \) is the distribution function of the \( i \)th particles, \( \kappa_i(x,p) \) is the Bohm-like parallel diffusion coefficient with a magnetic field strength \( B(x) \) and is given by
\[
\kappa_i(x,p) = \frac{u(x)}{3} r_i(x,p) = \frac{u(x)pc}{3ZeB(x)},
\]
(2)
\( Q_i(x,p) \) is the \( i \)th particle injection term and is given by
\[
Q_i(x,p) = \eta_i \frac{n_iu_0}{4\pi p_i^{3/2}} \delta (p - p_{i0}) \delta (x).
\]
(3)
where \( \eta_i \) is the fraction of the \( i \)th particles crossing the shock injected in the acceleration process, \( p_{i0} = Z_i p_{i0}\text{H} \) is the injection momentum of the \( i \)th particles, and \( \delta (x) \) is the position where particles are injected at the shock. In the thermal-leakage model [13, 20],
\[
p_{i0}\text{H} = \xi_u \sqrt{2m_i k_B T_{iH,2}},
\]
(4)
\[
T_{iH,2} = T_0 (R_{tot} / R_{atm})^{\gamma - 1} \left[ 1 + \frac{\gamma - 1}{\gamma + 1} \frac{p_{i0}}{p_{i0}\text{H}} \right],
\]
(5)
where \( k_B \) is the Boltzmann constant, \( m_i \) is proton mass, \( \gamma = 5/3 \) is the ratio of specific heats of gas and \( \xi_u \) is a parameter which defines \( p_{i0}\text{H} \) as a part of the momentum of the thermal protons in the downstream region.

In the absence of the dynamical back-reaction of the accelerated particles and magnetic field amplification at the shock, Eq. (13) can be independently solved. In other words, the pressure of the accelerated particles \( P_{CR}(x) \) and magnetic field pressure \( P_B(x) \) are negligible in comparison with the pressure \( P_i(x) \) of the gas with a density \( \rho \). In this case (called test particle approximation), the conservation equations of mass, momentum, and energy across the shock have the trivial solutions: \( \rho = \text{constant}, \ u = \text{constant}, \) and \( P_\text{sh} = \text{constant}. \)

In a general case, the two non-linear effects above must be considered.

(i) the dynamical back-reaction of the accelerated particles. Because the particles are efficiently accelerated at shock, the pressure \( P_{CR} \) of the accelerated particles must be included in the energy conservation equation, which results in the slow down of upstream plasma velocity in the shock frame, forming a so-called dynamical shock precursor. The normalized pressure of the accelerated particles can be estimated as
\[
P_{CR}(x) = \frac{4\pi}{3 \rho_0 u_0^5} \sum_{i} \int_{p_{i0}}^{p_{i0}\text{H}} dp \ p^3 \tilde{u}(x) f_i(x,p).
\]
(6)

(ii) the effect of the magnetic field amplification. Here magnetic field amplification due to the streaming instability of plasma flow is considered. In this case, the streaming instability of plasma flow will lead to the magnetic field amplification, and then the magnetic pressure \( (P_B) \) will have a significant change to affect shock compression ratios and particle’s spectra. Following Ref. [15], the normalized pressure of the magnetic field can be expressed as
\[
P_B(x) = \frac{2}{25} \frac{[1 - U(x)^{5/4}]^2}{U(x)^{3/2}},
\]
(7)
where \( U(x) = (\rho/\rho_0)[u(x)^2/u_0^2] \). And then the amplified magnetic field is estimated as
\[
B(x) = \sqrt{8\pi P_B(x)}.
\]
(8)
Considering both of the effects mentioned above, the momentum conservation equation, normalized to \( \rho_0 u_0^5 \), can be represented as:
\[
U(x) + P_{CR}(x) + P_B(x) + P_i(x) = 1 + \frac{1}{\gamma M_i^2},
\]
(9)
where \( P_{CR}(x) \) and \( P_B(x) \) are given by Eqs. (4) and (7), respectively, and \( P_i(x) = U(x)^{\gamma - 1} \) is the normalized pressure of the background gas with adiabatic index \( \gamma \), \( M_i = \frac{u}{c_s} \) is the Mach number of the fluid at upstream infinity.

Equation (11) with the spatial boundary condition \( f_i(x,0,p) \to 0 \) has been solved and the solution can be expressed as [19]
\[
f_i(x,p) = f_{sh,i}(p) e^{-\int_0^x dx' \kappa_i(x',p) \left[ 1 - \frac{U_i(x,p)}{U(x,p)} \right]}.
\]
(10)
where \( f_{sh,i}(p) \) is the distribution function at shock, which is
\[
f_{sh,i}(p) = \frac{n_i u_0 q_{p,i}(p)}{4\pi p_i^{3/2}} \left\{ \int_{p_{i0}}^{p_{i0}\text{H}} dp' \eta_i(p') \left[ \left[ U_{p,i}(p') \right]^{\gamma - 1} \right] \right\},
\]
(11)
where \( U_{p,i}(p) = U_1 - \int_{q_{p,i}}^{p_{i0}} dx [dU(x)/dx][f_i(x,p)/f_{sh,i}(p)] \) and \( q_{p,i}(p) = 3R_{tot} / (R_{tot} U_{p,i}(p) - 1) \); the function \( W_i(x,p) \) in Eq. (10) is given by
\[
W_i(x,p) = \int_0^x dx' \kappa_i(x',p) e^{-\int_0^{x'} dx'' \kappa_i(x'',p)} \left[ 1 - \frac{U_i(x',p)}{U(x',p)} \right],
\]
(12)
and \( W_{i,0}(p) = W_i(p)|_{x=x_0} \). The escape flux can be written by
\[
\Phi_{esc,i}(p) = -\kappa_i(x,p) \left. \frac{\partial f_i}{\partial x} \right|_{x_0} = -\frac{u_0 f_{sh,i}(p)}{W_{i,0}(p)}.
\]
(13)
Therefore, Eq. (11) is coupled with the equations describing mass, momentum, and energy flux conservations, leading to the problem of NLDSA whose solution can be obtained through the iterative method described in Refs. [19] and [21].
3 Particle Injection from a SNR

To perform our calculation, the following assumptions about a SNR’s evolution are made: the SNR is produced in a supernova explosion with an energy \( E_{\text{SN}} = 10^{51} \text{ erg} \) and an ejecta mass \( M_{\text{ej}} = 1.4 M_\odot \), and a shock moves with velocity \( u_0 = 4000 \text{ km/s} \) in a homogeneous and hot medium with a number density \( n_0 = 0.01 \text{ cm}^{-3} \), a temperature \( T_0 = 10^6 \text{ K} \) and a background magnetic field \( B_0 = 5 \mu \text{G} \). According to the analytical recipe given in Ref. [22], the SNR evolution is divided into two stages [12]:

1. Ejecta-dominated stage with \( \tau = t/T_{\text{ST}} \leq 1 \), both radius and velocity of the SNR are given by

   \[
   R_{\text{sh}}(t) \approx 14.1 \tau^{4/7} \text{ pc}, \quad V_{\text{sh}}(t) \approx 4140 \tau^{-3/7} \text{ km/s},
   \]

   where \( T_{\text{ST}} \approx 2000 \text{ yr} \) is used;

2. Sedov - Taylor stage with \( \tau = t/T_{\text{ST}} \geq 1 \),

   \[
   R_{\text{sh}}(t) \approx 16.2(\tau - 0.3)^{2/5} \text{ pc},
   \]

   \[
   V_{\text{sh}}(t) \approx 3330(\tau - 0.3)^{-3/5} \text{ km/s}.
   \]

Adiabatic loss is included because of the shell expansion. The energy \( E(t) \) of a particle with energy \( E_0 \) advected downstream at time \( t_0 \) is given by [12] where \( 4/3 \leq \gamma \leq 5/3 \).

The accelerated particles can escape from a shell volume \( V = 4\pi R_{\text{sh}}^2 dR_{\text{sh}} \), where \( dR_{\text{sh}} = V_{\text{sh}}(t)dt = [V_{\text{sh}}(t)/r_{\text{tot}}]dt \) and \( dt \) is the time increment. Since the particle numbers per unit volume per unit energy can be expressed as \( J_i(E, t) = 4\pi \rho_0 f_{\text{sh},i}(p) dp/dE \), where \( f_{\text{sh},i}(p) \) is the distribution function at shock radius \( R_{\text{sh}}(t) \), the particle numbers per unit energy in the shell volume \( V \) can be estimated as \( J_i(E, t) \times V \). Because the shell volume evolves with time \( t \) during SNR evolution, the particle numbers per unit energy are

\[
\phi_i(E) = 4\pi \int_{T_i}^{T_f} J_i[E, t] R_{\text{sh}}^2(t)[V_{\text{sh}}(t)/r_{\text{tot}}]dT \tag{18}
\]

where \( T_i \) and \( T_f \) are the initial and final times of the SNR evolution, here \( T_i = 0.1 T_{\text{ST}} \) and \( T_f = 15 T_{\text{ST}} \) with \( T_{\text{ST}} = 2000 \text{ yr} \) are used. There are three kinds of components for the particles escaping from the SNR [12]:

(i) The numbers per unit energy of particles which instantaneously escape around a maximum momentum \( p_{\text{max}}(t) \) from the upstream free escape boundary at \( x = x_0 \) during the Sedov - Taylor (ST) stage, where \( p_{\text{max}}(t) \) is determined by the finite size of the SNR during the ST stage, which can be estimated by \( \kappa_i(p_{\text{max}})/V_{\text{sh}}(t) \approx \chi R_{\text{sh}}(t), \) assuming \( x_0 \) to be a fraction \( \chi \) of the radius \( R_{\text{sh}}(t) \) of the SNR with a shock velocity \( V_{\text{sh}}(t) \). In this case, the number density of particles per unit energy is \( 4\pi \rho_0 f_{\text{sh},i}(p)/V_{\text{sh}}(t) \), where \( f_{\text{sh},i}(p) \) is given by Eq. (13), the particle numbers per unit energy are

\[
q_{\text{esc},i}(E_o) = \frac{16\pi^2}{c^2} \int_{T_i}^{T_f} p R_{\text{sh}}^2(t) E_o \Phi_{\text{esc},i}(p) dt. \tag{19}
\]

(ii) The numbers per unit energy of particles which are advected in the downstream region, leading to adiabatic losses as a consequence of the shell expansion, where the particles can escape at \( p > p_{\text{esc}}(t) \) at any given time and \( p_{\text{esc}}(t) \) can be estimated by \( \kappa_i(p_{\text{esc}}, B_2)/V_2 = x_0 \) with \( V_2 = V_{\text{sh}}/r_{\text{tot}} \). In this case, the particle numbers per unit energy are

\[
q_{\text{adv},i}(E_o) = \frac{16\pi^2}{c^2} \int_{T_i}^{T_f} p R_{\text{sh}}^2(t) E_o f_{\text{sh},i}(p) \frac{V_{\text{sh}}(t)}{r_{\text{tot}}} \times \left( \frac{V_{\text{sh}}(t)}{V_{\text{sh}}(t_0)} \right)^{2/3\gamma} dt. \tag{20}
\]

(iii) The numbers per unit energy of particles which escape the acceleration region from a broken shell at the end of a SNR’s evolution,

\[
q_{\text{shell},i}(E_o) = \lambda \times \frac{16\pi^2}{c^2} \int_{T_i}^{T_f} p R_{\text{sh}}^2(t) E_o f_{\text{sh},i}(p) \frac{V_{\text{sh}}(t)}{r_{\text{tot}}} dt, \tag{21}
\]

where the fraction of downstream escaping particles is taken as \( \lambda \approx 10\% \).

Therefore, the numbers per unit energy of the \( i \)th particles escaping from a single SNR is the sum of Eqs. (19) - (21), i.e.,

\[
q_i(E) = q_{\text{esc},i}(E) + q_{\text{adv},i}(E) + q_{\text{shell},i}(E). \tag{22}
\]

As an example, the spectra of H and He nuclei injected into the interstellar space are shown in Fig. 1. The model parameters are as follows: \( \xi_H = 3.0, \chi = 0.5 \quad T_0 = 10^6 \text{ K}, \quad n_0 = 0.01 \text{ cm}^{-3}, \) and \( B_0 = 5 \mu \text{G} \). In this figure, the spectra of three components mentioned above for H and He nuclei have been shown. From the figure, the first component dominates high-energy end, and the second component is a main contributor at lower energy. Moreover, from the bell-shaped curves (dotted line) which the particles have escaped to SNR from the upstream boundary, we can see that escaping occurs at highest energies. On the other hand, the cut-off energy of He is larger than that of H. For parameters used here, the shock is modified by the accelerated particles: \( r_{\text{sub}} = 3.41 \) and \( r_{\text{tot}} = 4.73 \).
Note that one parameter has an important influence on the CR spectrum, i.e., the parameter $\xi_H$ which describes the acceleration efficiency. In Fig. 2 it can be seen that the lower acceleration efficiency (i.e., a larger $\xi_H$) the particles, the flatter the resulting spectra and smaller the maximum energy.

4 Proton and He spectra Observed at the Earth

In this section, the spectra of H and He nuclei observed at the Earth are calculated. It is assumed that the propagation can be approximated by a simple leaky box model with 3 percent of SN explosion rate in our Galaxy. In this case, the energy spectrum $N_i(E)$ of the $i$th particles observed at the Earth is given by

$$N_i(E) \propto q_i(E) \left( \frac{1}{\lambda_{\text{esc},i}} + \frac{1}{\lambda_{\text{int},i}} \right)^{-1}, \quad (23)$$

where the escape path length $\lambda_{\text{esc},i}$ is a function of the particle magnetic rigidity $R_i = p/\beta Z_i$ and is approximated as $\lambda_{\text{esc}} = 7.3(R_i/10 \text{ GeV})^{-\delta} \beta(p)$ g/cm$^2$, where $\delta = 0.3 - 0.6$, $\beta(p)$ is the dimensionless speed of a nucleus of momentum $p$, $\lambda_{\text{int},i}$ is the interaction length and $\lambda_{\text{int},i} = \lambda_{0,i}(E/10 \text{ GeV})^{-\epsilon_i}$, where $\lambda_{0,H} = \lambda_{0,He} = 50$ g cm$^{-2}$, $\epsilon_H = 0.05$, $\epsilon_{He} = 0.0416$.

In order to estimate the value of $\eta_i$ from the measurement, following Ref. [24], the ratio of abundances between ions and protons at the same momentum $p = 10^5$ GeV/c measured at the Earth is defined as $K_{iH} = n_i/n_H$, and $\eta_i/\eta_H \approx K_{iH}Z_i^{-(\beta+\delta-3)}$. The distribution function is assumed to be a power law with a slope $\beta$, which can be obtained in the test particle approximation. Here $\beta + \delta = 4.7$ is used [24], so $\eta_{He}/\eta_H \approx 0.31K_{HeH}$.

Figure 3 shows the comparison of our model results with the H + He flux observed by ARGO-YBJ experiments [7, 23, 24]. The H + He spectrum has a knee feature of $\sim 700$ TeV [3]. For comparison, the observed data from various experiment groups are also shown. It can be seen from the figure that our model results can reproduce the flux observed by ARGO-YBJ experiments well. Moreover, the knee feature can be explained by the sum of H and He spectra but He spectrum dominates the high energy end. Note that the acceleration is efficient in this case, $\xi_H = 3.8$ corresponds to $\eta_H \approx 6.5 \times 10^{-5}$ and $\eta_{He} \approx 0.31K_{HeH} \times \eta_H \sim 2.01 \times 10^{-5}$. And the deviation at low energy may be due to solar modulation, and we did not consider the factor in our model.
As mentioned above, the acceleration efficiency plays an important role for the cut-off energy of the spectrum in the model. KASCADE experiments\cite{16} show that the observed H spectrum has a knee feature at a few PeV, which is not consistent with that observed by ARGO-YBJ experiment\cite{7, 25, 26}. To reproduce the observed H flux by KASCADE experiments with our model, the acceleration efficiency (i.e. $\xi_H$) is properly adjusted but other parameters are not changed for the proton injection spectrum, and the comparison of model result with the observed H flux is shown in Fig. 3. With $\xi_H = 3.5$ and $\delta = 0.5$, the observed H flux can be reproduced well in this model.

5 Results and Discussion

In this paper, the observed spectra of the light component (P+He) by ARGO-YBJ and the proton spectrum by KASCADE experiments are reproduced in the frame of non-linear diffusive shock acceleration model of the SNR. In the model, the escape spectrum of i-th particles injected into the ISM consists of three components (see Eqs. (11), (20) and (21)). In our calculation, for the case with ARGO-YBJ, parameter $\xi_H = 3.8$ is used, which corresponds to $\eta_H \approx 6.5 \times 10^{-5}$ and $\eta_{He} \approx 0.31 K_{He} H \times \eta_H \sim 2.01 \times 10^{-5}$, and $\eta_{He} \approx 4.32 \times 10^{-4}$ for the case with KASCADE. These values indicate that the acceleration at the SNR shock is very inefficient. In fact, it is generally believed that the acceleration is inefficient when $\eta_H \lesssim 10^{-5}$, which implies the acceleration efficiency is less than a few $\%$\cite{15}.

Finally, it should be pointed out that a typical evolution scenario of SNRs in our Galaxy and a simple leaky box approximation of CR propagation are used in our calculations. In fact, the case is very complicated\cite{24} and the detailed processes of CR spectrum reaching the knee through NLDSA remain uncertain. Meanwhile, the variations of parameters have great influence on the distribution of the knee. The contributions of different classes of SNRs to the CR spectrum should be taken into account\cite{32} and more relativistic CR propagation model\cite{33} should be used. It can be predicted that the data with smaller uncertainty will constrain the parameters (e.g. the amplified magnetic field) and demonstrate the effectiveness of the model.

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